

ASYMPTOTICS FOR THE WAVE EQUATION ON DIFFERENTIAL FORMS ON KERR-DE SITTER SPACE

PETER HINTZ AND ANDRÁS VASY

ABSTRACT. We study asymptotics for solutions of Maxwell's equations, in fact of the Hodge-de Rham equation $(d + \delta)u = 0$ without restriction on the form degree, on a geometric class of stationary spacetimes with a warped product type structure (without any symmetry assumptions), which in particular include Schwarzschild-de Sitter spaces of all spacetime dimensions $n \geq 4$. We prove that solutions decay exponentially to 0 or to stationary states in every form degree, and give an interpretation of the stationary states in terms of cohomological information of the spacetime. We also study the wave equation on differential forms and in particular prove analogous results on Schwarzschild-de Sitter spacetimes. We demonstrate the stability of our analysis and deduce asymptotics and decay for solutions of Maxwell's equations, the Hodge-de Rham equation and the wave equation on differential forms on Kerr-de Sitter spacetimes with small angular momentum.

1. INTRODUCTION

Maxwell's equations describe the dynamics of the electromagnetic field on a 4-dimensional spacetime (M, g) . Writing them in the form $(d + \delta_g)F = 0$, where δ_g is the codifferential, for the electromagnetic field F (a 2-form) suggests studying the operator $d + \delta_g$, whose square $\square_g = (d + \delta_g)^2$ is the Hodge d'Alembertian, i.e. the wave operator on differential forms. It is then very natural to study solutions of $(d + \delta_g)u = 0$ or $\square_g u = 0$ without restrictions on the form degree. Important examples of spacetimes that fit into the class of spacetimes studied in the present paper are Schwarzschild-de Sitter spacetimes with spacetime dimension ≥ 4 , and by very simple perturbation arguments, we can readily analyze waves on perturbations of these, in particular on Kerr-de Sitter spaces. Concretely, a special case of our general results is:

Theorem 1. *Let (M, g_a) denote a non-degenerate Kerr-de Sitter space with black hole mass $M_\bullet > 0$, cosmological constant $\Lambda > 0$ and angular momentum a ,¹ more precisely a suitable neighborhood of the domain of outer communications, and denote by t_* a smooth time coordinate. Suppose $u \in C^\infty(M; \Lambda M)$ is a solution of the equation*

$$(d + \delta_{g_a})u = 0,$$

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¹The non-degeneracy requirement ensures that the cosmological horizon lies outside the event horizon. For example, if M_\bullet is fixed and $\Lambda > 0$ is small, sufficiently small values of a work.

with smooth initial data, and denote by u_j the form degree j part of u , $j = 0, \dots, 4$. Then u_2 decays exponentially in t_* to a stationary state, which is a linear combination of the t_* -independent 2-forms $u_{a,1}, u_{a,2}$. In the standard (Boyer-Lindquist) local coordinate system on Kerr-de Sitter space, $u_{a,1}$ and $u_{a,2}$ have explicit closed form expressions; in particular, on Schwarzschild-de Sitter space, $u_{0,1} = r^{-2} dt \wedge dr$, and $u_{0,2} = \omega$ is the volume element of the round unit 2-sphere. Moreover, u_1 and u_3 decay exponentially to 0, while u_0 decays exponentially to a constant, and u_4 to a constant multiple of the volume form.

Suppose now $u \in C^\infty(M; \Lambda M)$ instead solves the wave equation

$$\square_{g_a} u = 0$$

with smooth initial data, then the same decay as before holds for u_0, u_2 and u_4 , while u_1 decays exponentially to a member of a 2-dimensional family of stationary states, likewise for u_3 .

The Schwarzschild-de Sitter case of this theorem, i.e. the special case $a = 0$, will be proved in Section 4.2, and we give explicit expressions for all stationary states, see Theorems 4.3 and 4.5, and Section 5 provides the perturbation arguments, see in particular Theorem 5.1. For the explicit form of $u_{a,1}$ and $u_{a,2}$, see Remark 5.4. Notice that asymptotics and decay of differential form solutions to the wave equation are much stronger statements than corresponding statements for Maxwell's equations or for the Hodge-de Rham equation.

We stress that the main feature of the spacetimes (M, g) considered in this paper is a warped product type structure of the metric, whereas *we do not make any symmetry assumptions on M* . From a geometric point of view then, the main novelty of this paper is a general cohomological interpretation of stationary states, which in the above theorem are merely explicitly given. On a technical level, we show how to analyze quasinormal modes (also called resonances, further discussed below) for equations on vector bundles whose natural inner product is not positive definite. To stress the generality of the method, we point out that symmetries only become relevant in explicit calculations for specific examples such as Schwarzschild-de Sitter and Kerr-de Sitter spaces. Even then, the perturbation analysis around Schwarzschild-de Sitter space works without restrictions on the perturbation; only for the explicit form of the space $\langle u_{a,1}, u_{a,2} \rangle$ of stationary states do we need the very specific form of the Kerr-de Sitter metric. Thus, combining the perturbation analysis with the non-linear framework developed by the authors in [30], we can immediately solve suitable *quasilinear wave equations on differential forms* on Kerr-de Sitter spacetimes; see Remark 5.3. To put this into context, part of the motivation for the present paper is the black hole stability problem, see the lecture notes by Dafermos and Rodnianski [14] for background on this, and we expect that the approach taken here will facilitate the linear part of the stability analysis, which, when accomplished, rather directly gives the non-linear result when combined with the non-linear analysis in [30].

1.1. Outline of the general result. Going back to the linear problem studied in this paper, we proceed to explain the general setup in more detail. Let \bar{X} be a compact orientable $(n - 1)$ -dimensional manifold with boundary $Y = \partial \bar{X}$ and interior $X = \bar{X}^\circ$, and let $M = \mathbb{R}_t \times X$, which is thus n -dimensional. Denote the connected components of Y , which are of dimension $(n - 2)$, by Y_i , for i in a finite

index set I . We assume that M is equipped with the metric

$$g = \alpha(x)^2 dt^2 - h(x, dx), \quad (1.1)$$

where h is a smooth Riemannian metric on \bar{X} (in particular, incomplete) and α is a boundary defining function of X , i.e. $\alpha \in \mathcal{C}^\infty(\bar{X})$, $\alpha = 0$ on Y , $\alpha > 0$ in X and $d\alpha|_Y \neq 0$. We moreover assume that every connected component Y_i of Y , $i \in I$, has a collar neighborhood $[0, \epsilon_i)_\alpha \times (Y_i)_y$ in which h takes the form

$$h = \tilde{\beta}_i(\alpha^2, y) d\alpha^2 + k_i(\alpha^2, y, dy) \quad (1.2)$$

with $\tilde{\beta}_i(0, y) \equiv \beta_i > 0$ constant along Y_i . In particular, h is an *even* metric in the sense of Guillarmou [26].² We change the smooth structure on \bar{X} to only include even functions of α , and show how one can then extend the metric g to a stationary metric (denoted \tilde{g} , but dropped from the notation in the sequel) on a bigger spacetime $\tilde{M} = \mathbb{R}_{t_*} \times \tilde{X}$, where t_* is a shifted time coordinate. Since the operator $d + \delta$ commutes with time translations, it is natural to consider the normal operator family

$$\tilde{\delta}(\sigma) + \tilde{\delta}(\sigma) = e^{it_*\sigma}(d + \delta)e^{-it_*\sigma}$$

acting on differential forms (valued in the form bundle of M) on a slice of constant t_* , identified with \tilde{X} ; the normal operator family $\tilde{\square}(\sigma)$ of \square is defined completely analogously. The proper way to view the normal operator family is as a family of operators on the boundary at infinity of a bordified version of \tilde{M} , where one introduces $\tau = e^{-t_*}$ as a new ‘time’ coordinate and adds $\tau = 0$, i.e. future infinity, to the manifold \tilde{M} . Thus, the normal operator family naturally lives on $\tau = 0$; see [52, §3] for details.

Since the Hodge d’Alembertian (and hence the normal operator family $\tilde{\square}(\sigma)$) has a scalar principal symbol, it can easily be shown to fit into the microlocal framework developed by Vasy [52]; we prove this in Section 2. In particular, $\tilde{\square}(\sigma)^{-1}$ is a meromorphic family of operators in $\sigma \in \mathbb{C}$, and under the assumption that the inverse family $\tilde{\square}(\sigma)^{-1}$ verifies suitable bounds as $|\operatorname{Re} \sigma| \rightarrow \infty$ and $\operatorname{Im} \sigma > -C$ (for $C > 0$ small), one can deduce exponential decay of solutions to $\square u = 0$, up to contributions from a finite dimensional space of *resonances*. Here, resonances are poles of $\tilde{\square}(\sigma)^{-1}$, and resonant states (for simple resonances) are elements of the kernel of $\tilde{\square}(\sigma)$ for a resonance σ . Therefore, proving wave decay and asymptotics is reduced to studying *high energy estimates*, which depend purely on geometric properties of the spacetime and will be further discussed below, and the location of resonances as well as the spaces of resonant states. (For instance, resonances in $\operatorname{Im} \sigma > 0$ correspond to exponentially growing solutions and hence are particularly undesirable when studying non-linear problems.) Our main theorem is then:

Theorem 2. *The only resonance of $d + \delta$ in $\operatorname{Im} \sigma \geq 0$ is $\sigma = 0$, and 0 is a simple resonance. Zero resonant states are smooth, and the space \mathcal{H} of these resonant states is equal to $\ker \tilde{d}(0) \cap \ker \tilde{\delta}(0)$. (In other words, resonant states, viewed as t_* -independent differential forms on \tilde{M} , are annihilated by d and δ .) Using the grading $\tilde{\mathcal{H}} = \bigoplus_{k=0}^n \tilde{\mathcal{H}}^k$ of $\tilde{\mathcal{H}}$ by form degrees, there is a canonical exact sequence*

$$0 \rightarrow H^k(\bar{X}) \oplus H^{k-1}(\bar{X}, \partial\bar{X}) \rightarrow \tilde{\mathcal{H}}^k \rightarrow H^{k-1}(\partial\bar{X}). \quad (1.3)$$

²Thus, as we will show, de Sitter and Schwarzschild-de Sitter spaces fit into this framework, whereas asymptotically flat spacetimes like Schwarzschild (or Kerr) do not.

Furthermore, the only resonance of \square in $\text{Im } \sigma \geq 0$ is $\sigma = 0$. Zero resonant states are smooth, and the space $\tilde{\mathcal{K}} = \bigoplus_{k=0}^n \tilde{\mathcal{K}}^k$ of these resonant states, graded by form degree and satisfying $\tilde{\mathcal{K}}^k \supset \tilde{\mathcal{H}}^k$, fits into the short exact sequence

$$0 \rightarrow H^k(\bar{X}) \oplus H^{k-1}(\bar{X}, \partial\bar{X}) \rightarrow \tilde{\mathcal{K}}^k \rightarrow H^{k-1}(\partial\bar{X}) \rightarrow 0. \quad (1.4)$$

Lastly, the Hodge star operator on \tilde{M} induces natural isomorphisms $\star: \tilde{\mathcal{H}}^k \xrightarrow{\cong} \tilde{\mathcal{H}}^{n-k}$ and $\star: \tilde{\mathcal{K}}^k \xrightarrow{\cong} \tilde{\mathcal{K}}^{n-k}$, $k = 0, \dots, n$.

See Theorem 3.20 for the full statement, including the precise definitions of the maps in the exact sequences. In fact, the various cohomology groups in (1.3) and (1.4) correspond to various types of resonant differential forms, namely forms which are square integrable on X with respect to a natural Riemannian inner product on forms on M (obtained by switching the sign in (1.1)), as well as ‘tangential’ and ‘normal’ forms in a decomposition $u = u_T + \alpha^{-1} dt \wedge u_N$ of the form bundle corresponding to the warped product structure of the metric. Roughly speaking, (1.4) encodes the fact that resonant states for which a certain boundary component vanishes are square integrable with respect to the natural Riemannian inner product on X and can be shown to canonically represent absolute (for tangential forms) or relative (for normal forms) de Rham cohomology of \bar{X} , while the aforementioned boundary component is a harmonic form on Y and can be specified freely for resonant states of \square . (Notice by contrast that the last map in the exact sequence (1.3) for $d + \delta$ is not necessarily surjective.)

The proof of Theorem 2 proceeds in several steps. First, we exclude resonances in $\text{Im } \sigma > 0$ in Section 3.1; the idea here is to relate the normal operator family of $d + \delta$ (a family of operators on the extended space \tilde{X}) to another normal operator family $\hat{d}(\sigma) + \hat{\delta}(\sigma) = e^{it\sigma}(d + \delta)e^{-it\sigma}$, which is a family of operators on X that degenerates at $\partial\bar{X}$, but has the advantage of having a simple form in view of the warped product type structure (1.1) of the metric: Since one formally obtains $\hat{d}(\sigma) + \hat{\delta}(\sigma)$ by replacing each ∂_t in the expression for $d + \delta$ by $-i\sigma$, we see that on a formal level $\hat{d}(\sigma) + \hat{\delta}(\sigma)$ for purely imaginary σ resembles the normal operator family of the Hodge-de Rham operator of the Riemannian metric on M mentioned above; then one can show the triviality of $\ker(\hat{d}(\sigma) + \hat{\delta}(\sigma))$ in a way that is very similar to how one would show the triviality of $\ker(A + \sigma)$ for self-adjoint A and $\text{Im } \sigma > 0$. For not purely imaginary σ , but still with $\text{Im } \sigma > 0$, one can change the tangential part of the metric on M in (1.1) by a complex phase and then run a similar argument, using that the resulting ‘inner product’, while complex, still has some positivity properties. Next, in Section 3.2, we exclude non-zero real resonances by means of a boundary pairing argument, which is a standard technique in scattering theory, see e.g. Melrose [41]. Finally, the analysis of the zero resonance in Section 3.3 relies on a boundary pairing type argument, and we again use the Riemannian inner product on forms on M . The fact that this Riemannian inner product is singular at $\partial\bar{X}$ implies that resonant states are not necessarily square integrable, and whether or not a state is square integrable is determined by the absence of a certain boundary component of the state. This is a crucial element of the cohomological interpretation of resonant states in Section 3.4.

As already alluded to, deducing wave expansions and decay from Theorem 2 requires high energy estimates for the normal operator family. These are easy to obtain if the metric h on X is *non-trapping*, i.e. all geodesics escape to $\partial\bar{X}$, as is

the case for the static patch of de Sitter space (discussed in Section 4.1). Another instance in which suitable estimates hold is when the only trapping within X is *normally hyperbolic trapping*, as is the case for Kerr-de Sitter spaces with parameters in a certain range. In the scalar setting, such estimates are now widely available, see for instance Wunsch and Zworski [56], Dyatlov [23] and their use in Vasy [52]: The proof of exponential decay relies on high energy estimates in a strip below the real line. For \square acting on differential forms, obtaining high energy estimates requires a smallness assumption on the imaginary part of the subprincipal symbol of \square relative to a *positive definite* inner product on the form bundle; the choice of inner product affects the size of the subprincipal symbol. Conceptually, the natural framework in which to find such an inner product involves *pseudodifferential inner products*. This notion was introduced by Hintz [28] and used there to prove high energy estimates for \square on tensors of arbitrary rank on perturbations of Schwarzschild-de Sitter space. In the present paper, we use the estimates provided in [28] as black boxes.

1.2. Previous and related work. The present paper seems to be the first to describe asymptotics for differential forms solving the wave or Hodge-de Rham equation in all form degrees and in this generality, and also the first to demonstrate the forward solvability of non-scalar quasilinear wave equations on black hole spacetimes, but we point out that for applications in general relativity, our results require the cosmological constant to be positive, whereas previous works on Maxwell's equations deal with asymptotically flat spacetimes: Sterbenz and Tataru [46] showed local energy decay for Maxwell's equations on a class of spherically symmetric asymptotically flat spacetimes including Schwarzschild.³ Blue [6] established conformal energy and pointwise decay estimates in the exterior of the Schwarzschild black hole; Andersson and Blue [2] proved similar estimates on slowly rotating Kerr spacetimes. These followed earlier results for Schwarzschild by Inglese and Nicolo [32] on energy and pointwise bounds for integer spin fields in the far exterior of the Schwarzschild black hole, and by Bachelot [3], who proved scattering for electromagnetic perturbations. Finster, Kamran, Smoller and Yau [24] proved local pointwise decay for Dirac waves on Kerr. There are further works which in particular establish bounds for certain components of the Maxwell field, see Donniger, Schlag and Soffer [18] and Whiting [55]. Dafermos [10], [11] studied the non-linear Einstein-Maxwell-scalar field system under the assumption of spherical symmetry.

Vasy's proof of the meromorphy of the (modified) resolvent of the Laplacian on differential forms on asymptotically hyperbolic spaces [51] makes use of the same microlocal framework as the present paper, and it also shows how to link the 'intrinsic' structure of the asymptotically hyperbolic space and the form of the Hodge-Laplacian with a 'non-degenerately extended' space and operator. For Kerr-de Sitter spacetimes, Dyatlov [20] defined quasinormal modes or resonances in the same way as they are used here, and obtained exponential decay to constants away from the event horizon for scalar waves. This followed work of Melrose, Sá Barreto and Vasy [40], where this was shown up to the event horizon of a Schwarzschild-de Sitter black hole, and of Dafermos and Rodnianski [13] who proved polynomial decay in this setting. Dyatlov proved exponential decay up to the event horizon for Kerr-de Sitter in [19], and significantly strengthened this in [21], obtaining a full

³One needs to assume the vanishing of the electric and the magnetic charge. For a positive cosmological constant, this precisely corresponds to assuming the absence of the $r^{-2} dt \wedge dr$ and ω asymptotics in form degree 2 in the Schwarzschild-de Sitter case of Theorem 1.

resonance expansion for scalar waves, improving on the result of Bony and Häfner [7] in the Schwarzschild-de Sitter setting, which in turn followed Sá Barreto and Zworski [4].

In the scalar setting too, the wave equation on asymptotically flat spacetimes has received more attention. Dafermos, Rodnianski and Shlapentokh-Rothman [17], building on [15, 16, 44], established the decay of scalar waves on all non-extremal Kerr spacetimes, following pioneering work by Kay and Wald [33, 54] in the Schwarzschild setting. Tataru and Tohaneanu [47, 48] proved decay and Price's law for slowly rotating Kerr using local energy decay estimates, and Strichartz estimates were proved by Marzuola, Metcalfe, Tataru and Tohaneanu [37].

Non-linear results for wave equations on black hole spacetimes include [30], see also the references therein, Luk's work [36] on semilinear forward problems on Kerr, and the scattering construction of dynamical black holes by Dafermos, Holzegel and Rodnianski [12]. Fully general stability results for Einstein's equations specifically are available for de Sitter space by the works of Friedrich [25], Anderson [1], Rodnianski and Speck [43] and Ringström [42], and for Minkowski space by the work of Christodoulou and Klainerman [9], partially simplified and extended by Lindblad and Rodnianski [34, 35], Bieri and Zipser [5] and Speck [45].

1.3. Structure of the paper. In Section 2, we show how to put the Hodge-de Rham and wave equation on differential forms into the microlocal framework of [52]; this is used in Section 3 in the analysis of resonances in $\text{Im } \sigma \geq 0$, and we prove Theorem 2 there. In Section 4, we apply this result on de Sitter space, where we can take a global point of view which simplifies explicit calculations considerably, and on Schwarzschild-de Sitter space, where such a global picture is not available, but using Theorem 2, the necessary computations are still very straightforward. In Section 5, we show the perturbation stability of the analysis, in particular deal with Kerr-de Sitter space, and indicate how this gives the forward solvability for quasilinear wave equations on differential forms.

2. ANALYTIC SETUP

Recall that we are working on a spacetime $M = \mathbb{R}_t \times X$, equipped with a metric g as in (1.1)-(1.2), where X is the interior of a compact orientable manifold \overline{X} with boundary $Y = \partial\overline{X}$ and boundary defining function $\alpha \in \mathcal{C}^\infty(\overline{X})$. Fixing a collar neighborhood of Y identified with $[0, \epsilon)_\alpha \times Y$, denote by $\overline{X}_{\text{even}}$ the manifold \overline{X} with the smooth structure changed so that only even functions in α are smooth, i.e. smooth functions are precisely those for which all odd terms in the Taylor expansion at all boundary components vanish. For brevity, we assume from now on that Y is connected,

$$h = \tilde{\beta}(\alpha^2, y)^2 d\alpha^2 + k(\alpha^2, y, dy) \quad (2.1)$$

in a collar neighborhood of Y , and thus $\tilde{\beta}(0, y) \equiv \beta$ is a single constant, but all of our arguments readily go through in the case of multiple boundary components. The main examples of spaces which directly fit into this setup are the static patch of de Sitter space (with 1 boundary component) and Schwarzschild-de Sitter space (with 2 boundary components); see Section 4 for details.

On M , we consider the Hodge-de Rham operator $d + \delta$, acting on differential forms. We put its square, the Hodge d'Alembertian $\square = (d + \delta)^2$, which is principally scalar, into the microlocal framework developed in [52]: We renormalize the

time coordinate t in the collar neighborhood of Y by writing

$$t = t_* + F(\alpha), \quad \partial_\alpha F(\alpha) = -\frac{\tilde{\beta}}{\alpha} - 2\alpha c(\alpha^2, y) \quad (2.2)$$

with c smooth, hence $F(\alpha) \in -\beta \log \alpha + \mathcal{C}^\infty(\overline{X}_{\text{even}})$; notice that the above requirement on F only makes sense near Y . We introduce the boundary defining function $\mu = \alpha^2$ of $\overline{X}_{\text{even}}$; then one computes

$$g = \mu dt_*^2 - (\tilde{\beta} + 2\mu c) dt_* d\mu + (\mu c^2 + \tilde{\beta} c) d\mu^2 - k(\mu, y, dy). \quad (2.3)$$

In particular, the determinant of g in these coordinates equals $-\frac{\tilde{\beta}^2}{4} \det(k)$, hence g is non-degenerate up to Y . Furthermore, we claim that we can choose $c(\mu, y)$ such that ∂_{t_*} is timelike on $\mathbb{R}_{t_*} \times \overline{X}_{\text{even}}$ (this requirement is explained below and in [52, §7]); indeed, with G denoting the dual metric to g , we require

$$G(\partial_{t_*}, \partial_{t_*}) = -4\tilde{\beta}^{-2}(\mu c^2 + \tilde{\beta} c) > 0. \quad (2.4)$$

This is trivially satisfied if $c = -\tilde{\beta}/2\mu$, which corresponds to undoing the change of coordinates in (2.2), however we want c to be smooth at $\mu = 0$. But for $\mu \geq 0$, (2.4) holds provided $-\tilde{\beta}/\mu < c < 0$; hence, we can choose a smooth c verifying (2.4) in $\mu \geq 0$ and such that moreover $c = -\tilde{\beta}/2\mu$ in $\mu \geq \mu_1$ (intersected with the collar neighborhood of Y) for any fixed $\mu_1 > 0$. Thus, we can choose F as in (2.2) with $F = 0$ in $\alpha^2 \geq \mu_1$ (in particular, F is defined globally on X) such that (2.4) holds.

Since the metric g in (2.3) is stationary (t_* -independent) and non-degenerate on $\overline{X}_{\text{even}}$, it can be extended to a stationary Lorentzian metric on an extension \tilde{X}_δ into which $\overline{X}_{\text{even}}$ embeds. Concretely, one defines $\tilde{X}_\delta = (\overline{X}_{\text{even}} \sqcup ([-\delta, \epsilon)_\mu \times Y_y)) / \sim$ with the obvious smooth structure, where \sim identifies elements of $[0, \epsilon)_\mu \times Y_y$ with points in $\overline{X}_{\text{even}}$ by means of the collar neighborhood of Y . Then, extending $\tilde{\beta}$ and k , and thus g , in an arbitrary t_* -independent manner to \tilde{X}_δ , the extended metric, which we denote by \tilde{g} , is non-degenerate on \tilde{X}_δ for sufficiently small $\delta > 0$, and ∂_{t_*} remains timelike uniformly on $\mathbb{R}_{t_*} \times \tilde{X}_\delta$: Indeed, in $\mu < 0$, (2.4) (with the dual metric \tilde{G} of \tilde{g} in place of G) holds for any negative function c as long as $\tilde{\beta}$ remains positive on \tilde{X}_δ . Reducing $\delta > 0$ further if necessary (to enforce the relevant structure of the null-geodesic flow near Y within $\tilde{X}_\delta \setminus \overline{X}_{\text{even}}$, see [52, §2]), we let \tilde{X} be the double space of \tilde{X}_δ , which is thus a compact manifold without boundary, and denote by \tilde{g} the extended metric on \tilde{X} , slightly modified near $\partial\tilde{X}_\delta$ to ensure the smoothness of \tilde{g} on the double space \tilde{X} .

The operator $d + \delta_g$ on M now extends to an operator $d + \delta_{\tilde{g}}$ on $\tilde{M} = \mathbb{R}_{t_*} \times \tilde{X}$. Correspondingly, the wave operator \square_g on M extends to the wave operator $\square_{\tilde{g}}$ on \tilde{M} . Denote the conjugation of $\square_{\tilde{g}}$ on \tilde{M} by the Fourier transform in $-t_*$ (equivalently, by the Mellin transform in e^{-t_*}) by $\tilde{\square}_{\tilde{g}}(\sigma)$, that is to say (dropping the subscript \tilde{g} for brevity)

$$\tilde{\square}(\sigma) = e^{it_*\sigma} \square e^{-it_*\sigma};$$

since \square is invariant under translations in t_* , this amounts to replacing each ∂_{t_*} in the expression for \square by $-i\sigma$. The operator $\tilde{\square}(\sigma)$ acts on sections of the pullback $\Lambda_{\tilde{X}} \tilde{M}$ of the form bundle $\Lambda \tilde{M}$ under the map $\tilde{X} \rightarrow \tilde{M}$, $\tilde{x} \mapsto (0, \tilde{x})$, and writing

differential forms \tilde{u} on \tilde{M} as

$$\tilde{u} = \tilde{u}_T + dt_* \wedge \tilde{u}_N \quad (2.5)$$

with \tilde{u}_T and \tilde{u}_N valued in forms on \tilde{X} , we can identify $\Lambda_{\tilde{X}} \tilde{M}$ with $\Lambda \tilde{X} \oplus \Lambda \tilde{X}$.

The last step required to show that \square , more precisely $\square(\sigma)$, fits into the framework described in [52] is classical non-trapping for the bicharacteristic flow of $\square(\sigma)$; complex absorption can be dealt with by the arguments of [52, §§3-4]. But in fact, we even have ellipticity in X : Indeed, on X , we have

$$\tilde{\square}(\sigma) = e^{-iF\sigma} e^{it\sigma} \square e^{-it\sigma} e^{iF\sigma} = e^{-iF\sigma} \hat{\square}(\sigma) e^{iF\sigma}, \quad (2.6)$$

where $\hat{\square}(\sigma) = e^{it\sigma} \square e^{-it\sigma}$ is the conjugation of \square by the Fourier transform in $-t$, and F is as in (2.2); here, we view $\hat{\square}(\sigma)$ as an operator acting on sections of $\Lambda_{\tilde{X}} \tilde{M}|_X$. Now, the latter bundle can be identified with $\Lambda X \oplus \Lambda X$ by writing differential forms as $u = u_T + dt \wedge u_N$, with u_T and u_N valued in forms on X , and switching between this identification and (2.5) amounts to conjugating $\hat{\square}(\sigma)$ by a bundle isomorphism on $\Lambda X \oplus \Lambda X$, which preserves ellipticity. The standard principal symbol of $\hat{\square}(\sigma)$ as a second order operator acting on sections of $\Lambda X \oplus \Lambda X$ is given by $(-H) \oplus (-H)$, where H is the dual metric to h , here identified with the dual metric function on T^*X ; this follows from the calculations in the next section. Since H is Riemannian, this implies that $\hat{\square}(\sigma)$, hence $\tilde{\square}(\sigma)$, is classically elliptic in X , which trivially implies the non-trapping property.

Hence by [52, Theorem 7.3], $\tilde{\square}(\sigma)$ is an analytic family of Fredholm operators on suitable function spaces, and the inverse family $\tilde{\square}(\sigma)^{-1} : \mathcal{C}^\infty(\tilde{X}; \Lambda \tilde{X} \oplus \Lambda \tilde{X}) \rightarrow \mathcal{C}^{-\infty}(\tilde{X}; \Lambda \tilde{X} \oplus \Lambda \tilde{X})$ (where we use the identification (2.5)) admits a meromorphic continuation from $\text{Im } \sigma \gg 0$ to the complex plane; note however that without further assumptions on the geodesic flow (for instance, semiclassical non-trapping or normally hyperbolic trapping), we do not obtain any high energy bounds. Moreover (see [52, Lemma 3.5]), the Laurent coefficient at the poles are finite rank operators mapping sufficiently regular distributions to elements of $\mathcal{C}^\infty(\tilde{X}; \Lambda \tilde{X} \oplus \Lambda \tilde{X})$.

For present purposes, it is actually more convenient to replace complex absorption by Cauchy hypersurfaces outside of \bar{X}_{even} as in [29, §2] and [27, §8]; the above properties on $\tilde{\square}(\sigma)^{-1}$ hold true in this setting as well. We then deduce:

Lemma 2.1. *A complex number $\sigma \in \mathbb{C}$ is a resonance of \square , i.e. $\tilde{\square}(\sigma)^{-1}$ has a pole at σ , if and only if there exists a non-zero $u \in \alpha^{-i\beta\sigma} \mathcal{C}^\infty(\bar{X}_{\text{even}}; \Lambda \bar{X}_{\text{even}} \oplus \Lambda \bar{X}_{\text{even}})$ (using the identification (2.5)) such that $\hat{\square}(\sigma)u = 0$.*

Proof. If $\sigma \in \mathbb{C}$ is a resonance, then there exists a non-zero $\tilde{u} \in \mathcal{C}^\infty(\tilde{X}; \Lambda \tilde{X} \oplus \Lambda \tilde{X})$ with $\tilde{\square}(\sigma)\tilde{u} = 0$. Restricting to \bar{X} , this implies by (2.6) and (2.2) that $\hat{\square}(\sigma)u = 0$ for $u = e^{iF\sigma} \tilde{u}|_{\bar{X}} \in \alpha^{-i\beta\sigma} \mathcal{C}^\infty(\bar{X}_{\text{even}}; \Lambda \bar{X}_{\text{even}} \oplus \Lambda \bar{X}_{\text{even}})$. If $u = 0$, then \tilde{u} vanishes to infinite order at Y , and since $\tilde{\square}(\sigma)$ is a conjugate of a wave or Klein-Gordon operator on an asymptotically de Sitter space, see [53], unique continuation at infinity on the de Sitter side as in [50, Proposition 5.3] (which is in the scalar setting, but works similarly in the present context since it relies on a semiclassical argument in which only the principal symbol of the wave operator matters, and this is the same in our setting) shows that $\tilde{u} \equiv 0$ on \tilde{X} .⁴ Hence, $u \neq 0$, as desired.

⁴This is the place where we use that we capped off \tilde{X} outside of \bar{X}_{even} by a Cauchy hypersurface.

Conversely, given a $u \in \alpha^{-i\beta\sigma} \mathcal{C}^\infty(\overline{X}_{\text{even}}; \Lambda \overline{X}_{\text{even}} \oplus \Lambda \overline{X}_{\text{even}})$ with $\widehat{\square}(\sigma)u = 0$, we define $\tilde{u}' \in \mathcal{C}^\infty(\tilde{X}; \Lambda \tilde{X} \oplus \Lambda \tilde{X})$ to be any smooth extension of $e^{-iF\sigma}u$ from $\overline{X}_{\text{even}}$ to \tilde{X} . Then $\tilde{\square}(\sigma)\tilde{u}'$ is identically zero in X and thus vanishes to infinite order at Y ; hence, we can solve

$$\tilde{\square}(\sigma)\tilde{v} = -\tilde{\square}(\sigma)\tilde{u}'$$

in $\tilde{X} \setminus X$ with \tilde{v} vanishing to infinite order at Y ; thus, extending \tilde{v} by 0 to X , we find that $\tilde{u} = \tilde{u}' + \tilde{v}$ is a non-zero solution to $\tilde{\square}(\sigma)\tilde{u} = 0$ on \tilde{X} . \square

Since $\square = (d + \delta)^2$, we readily obtain the following analogue of Lemma 2.1 for $d + \delta$:

Lemma 2.2. *The map⁵ $\ker_{\mathcal{C}^\infty(\tilde{X})}(\tilde{d}(\sigma) + \tilde{\delta}(\sigma)) \rightarrow \ker_{\alpha^{-i\beta\sigma} \mathcal{C}^\infty(\overline{X}_{\text{even}})}(\widehat{d}(\sigma) + \widehat{\delta}(\sigma))$, $\tilde{u} \mapsto e^{iF\sigma}\tilde{u}|_X$, is an isomorphism.*

Proof. Since $\tilde{u} \in \ker(\tilde{d}(\sigma) + \tilde{\delta}(\sigma))$ implies $\tilde{u} \in \ker \tilde{\square}(\sigma)$, injectivity follows from the proof of Lemma 2.1. To show surjectivity, take $u \in e^{iF\sigma} \mathcal{C}^\infty(\overline{X}_{\text{even}})$ with $(\widehat{d}(\sigma) + \widehat{\delta}(\sigma))u = 0$ and choose any smooth extension \tilde{u}' of $e^{-iF\sigma}u$ to \tilde{X} . Solving $\tilde{\square}(\sigma)\tilde{v}' = -(\tilde{d}(\sigma) + \tilde{\delta}(\sigma))\tilde{u}'$ with $\text{supp } \tilde{v}' \subset \tilde{X} \setminus X$ and then defining $\tilde{v} = (\tilde{d}(\sigma) + \tilde{\delta}(\sigma))\tilde{v}'$, we see that $\tilde{u} = \tilde{u}' + \tilde{v}$ extends \tilde{u}' to \tilde{X} and is annihilated by $\tilde{d}(\sigma) + \tilde{\delta}(\sigma)$. \square

Thus, when studying the location and structure of resonances, we already have very precise information about regularity and asymptotics (on X) of potential resonant states.

Lastly, we remark that since $\tilde{\square}(\sigma) = (\tilde{d}(\sigma) + \tilde{\delta}(\sigma))^2$ is an analytic family of Fredholm operators with meromorphic inverse, the same holds for $\tilde{d}(\sigma) + \tilde{\delta}(\sigma)$. More precisely, $\tilde{\square}(\sigma): \mathcal{X}^s \rightarrow \mathcal{Y}^{s-1}$ with $\mathcal{X}^s = \{u \in H^s(\tilde{X}; \Lambda \tilde{X} \oplus \Lambda \tilde{X})^- : \tilde{\square}(\sigma)u \in H^{s-1}(\tilde{X}; \Lambda \tilde{X} \oplus \Lambda \tilde{X})\}$ and $\mathcal{Y}^{s-1} = H^{s-1}(\tilde{X}; \Lambda \tilde{X} \oplus \Lambda \tilde{X})^-$, with $(-)$ denoting extendible distributions [31, Appendix B], is Fredholm provided s is large enough (depending on $\text{Im } \sigma$), and thus $\tilde{d}(\sigma) + \tilde{\delta}(\sigma): \mathcal{Z}^s \rightarrow \mathcal{Y}^s$ with $\mathcal{Y}^s = \{u \in H^s(\tilde{X}; \Lambda \tilde{X} \oplus \Lambda \tilde{X})^- : (\tilde{d}(\sigma) + \tilde{\delta}(\sigma))u \in H^s(\tilde{X}; \Lambda \tilde{X} \oplus \Lambda \tilde{X})\}$ is Fredholm for the same s . In addition, $\tilde{d}(\sigma) + \tilde{\delta}(\sigma)$ acting on these spaces is invertible if and only if its square $\tilde{\square}(\sigma)$ is. In particular, $\tilde{d}(\sigma) + \tilde{\delta}(\sigma)$ has index 0, being an analytic family of Fredholm operators which is invertible for $\text{Im } \sigma \gg 0$.

3. RESONANCES IN $\text{Im } \sigma \geq 0$

Using Lemma 2.2, we now study the resonances of in $\text{Im } \sigma \geq 0$ by analyzing the operator $\widehat{d}(\sigma) + \widehat{\delta}(\sigma)$ (and related operators) on $\overline{X}_{\text{even}}$. Recall that a resonance at $\sigma \in \mathbb{C}$ and a corresponding resonant state \tilde{u} yield a solution $(d + \delta)(e^{-it_*\sigma}\tilde{u}) = 0$, hence $\text{Im } \sigma > 0$ implies in view of $|e^{-it_*\sigma}| = e^{t_* \text{Im } \sigma}$ that $e^{-it_*\sigma}\tilde{u}$ grows exponentially in t_* , whereas resonances with $\text{Im } \sigma = 0$ yield solutions which at most grow polynomially in t_* (and do not decay). *We will continue to drop the metric g or \tilde{g} from the notation for brevity.*

In order to keep track of fiber inner products and volume densities, we will use the following notation.

⁵We drop the bundles from the notation for simplicity.

Definition 3.1. For a density μ on X and a complex vector bundle $\mathcal{E} \rightarrow X$ equipped with a positive definite Hermitian form B , let $L^2(X, \mu; \mathcal{E}, B)$ be the space of all sections u of \mathcal{E} for which $\|u\|_{\mu, B}^2 := \int_X B(u, u) d\mu < \infty$.

If B is merely assumed to be sesquilinear (but not necessarily positive definite), we define the pairing

$$\langle u, v \rangle_{\mu, B} := \int_X B(u, v) d\mu$$

for all sections u, v of \mathcal{E} for which $B(u, v) \in L^1(X, \mu)$. If the choice of the density μ or inner product B is clear from the context, it will be dropped from the notation.

Remark 3.2. It will always be clear what bundle \mathcal{E} we are using at a given time, so \mathcal{E} will from now on be dropped from the notation; also, X will mostly be suppressed.

Since the metric g in (1.1) has a warped product structure and αdt has unit squared norm, it is natural to write differential forms on $M = \mathbb{R}_t \times X_x$ as

$$u(t, x) = u_T(t, x) + \alpha dt \wedge u_N(t, x), \quad (3.1)$$

where the tangential and normal forms u_T and u_N are t -dependent forms on X , and we will often write this as

$$u(t, x) = \begin{pmatrix} u_T(t, x) \\ u_N(t, x) \end{pmatrix}.$$

Thus, the differential d on M is given in terms of the differential d_X on X by

$$d = \begin{pmatrix} d_X & 0 \\ \alpha^{-1} \partial_t & -\alpha^{-1} d_X \alpha \end{pmatrix}. \quad (3.2)$$

Since the dual metric is given by $G = \alpha^{-2} \partial_t^2 - H$, the fiber inner product G_k on k -forms is given by

$$G_k = \begin{pmatrix} (-1)^k H_k & 0 \\ 0 & (-1)^{k-1} H_{k-1} \end{pmatrix}, \quad (3.3)$$

where H_q denotes the fiber inner product on q -forms on X . Furthermore, the volume density on M is $|dg| = \alpha |dt dh|$, and we therefore compute the $L^2(M, |dg|)$ -adjoint of d to be

$$\delta = \begin{pmatrix} -\alpha^{-1} \delta_X \alpha & -\alpha^{-1} \partial_t \\ 0 & \delta_X \end{pmatrix}, \quad (3.4)$$

where δ_X is the $L^2(X, |dh|; \Lambda X, H)$ -adjoint of d_X . Thus,

$$\widehat{d}(\sigma) = \begin{pmatrix} d_X & 0 \\ -i\sigma\alpha^{-1} & -\alpha^{-1} d_X \alpha \end{pmatrix}, \quad \widehat{\delta}(\sigma) = \begin{pmatrix} -\alpha^{-1} \delta_X \alpha & i\sigma\alpha^{-1} \\ 0 & \delta_X \end{pmatrix}. \quad (3.5)$$

In the course of our arguments we will need to justify various integrations by parts and boundary pairing arguments. This requires a precise understanding of the asymptotics of u_T and u_N for potential resonant states u at $Y = \partial \bar{X}_{\text{even}}$. To this end, we further decompose the bundle $\Lambda X \oplus \Lambda X$ near Y by writing u_T as

$$u_T = u_{TT} + d\alpha \wedge u_{TN} \quad (3.6)$$

and similarly for u_N , hence

$$u = u_{TT} + d\alpha \wedge u_{TN} + \alpha dt \wedge u_{NT} + \alpha dt \wedge d\alpha \wedge u_{NN}, \quad (3.7)$$

where the $u_{\bullet\bullet}$ are forms on X valued in ΛY . Now for a resonant state u , we have

$$u = \alpha^{-i\beta\sigma} (\tilde{u}'_{TT} + d(\alpha^2) \wedge \tilde{u}'_{TN} + dt_* \wedge \tilde{u}'_{NT} + dt_* \wedge d(\alpha^2) \wedge \tilde{u}'_{NN}) \quad (3.8)$$

near Y with $\tilde{u}'_{\bullet\bullet} \in \mathcal{C}^\infty(\overline{X}_{\text{even}}; \Lambda Y)$, which we rewrite in terms of the decomposition (3.7) using (2.2), obtaining

$$u = \alpha^{-i\beta\sigma} (\tilde{u}'_{TT} + d\alpha \wedge (2\alpha\tilde{u}'_{TN} - F'(\alpha)\tilde{u}'_{NT}) \\ + \alpha dt \wedge \alpha^{-1}\tilde{u}'_{NT} + 2\alpha dt \wedge d\alpha \wedge \tilde{u}'_{NN});$$

hence introducing the ‘change of basis’ matrix

$$\mathcal{C} = \begin{pmatrix} 1 & 0 & 0 & 0 \\ 0 & \alpha & \beta\alpha^{-1} & 0 \\ 0 & 0 & \alpha^{-1} & 0 \\ 0 & 0 & 0 & 1 \end{pmatrix}$$

and defining the space

$$\mathcal{C}_{(\sigma)}^\infty := \mathcal{C}\alpha^{-i\beta\sigma} \begin{pmatrix} \mathcal{C}^\infty(\overline{X}_{\text{even}}; \Lambda Y) \\ \mathcal{C}^\infty(\overline{X}_{\text{even}}; \Lambda Y) \\ \mathcal{C}^\infty(\overline{X}_{\text{even}}; \Lambda Y) \\ \mathcal{C}^\infty(\overline{X}_{\text{even}}; \Lambda Y) \end{pmatrix} \subset \begin{pmatrix} \alpha^{-i\beta\sigma}\mathcal{C}^\infty(\overline{X}_{\text{even}}; \Lambda Y) \\ \alpha^{-i\beta\sigma-1}\mathcal{C}^\infty(\overline{X}_{\text{even}}; \Lambda Y) \\ \alpha^{-i\beta\sigma-1}\mathcal{C}^\infty(\overline{X}_{\text{even}}; \Lambda Y) \\ \alpha^{-i\beta\sigma}\mathcal{C}^\infty(\overline{X}_{\text{even}}; \Lambda Y) \end{pmatrix}, \quad (3.9)$$

we obtain

$$\begin{pmatrix} u_{TT} \\ u_{TN} \\ u_{NT} \\ u_{NN} \end{pmatrix} = \mathcal{C}\alpha^{-i\beta\sigma} \begin{pmatrix} \tilde{u}_{TT} \\ \tilde{u}_{TN} \\ \tilde{u}_{NT} \\ \tilde{u}_{NN} \end{pmatrix} \in \mathcal{C}_{(\sigma)}^\infty \quad (3.10)$$

with $\tilde{u}_{\bullet\bullet} \in \mathcal{C}^\infty(\overline{X}_{\text{even}}; \Lambda Y)$, where the $u_{\bullet\bullet}$ are the components of u in the decomposition (3.7).

We will also need the precise form of $\hat{d}(\sigma)$ and $\hat{\delta}(\sigma)$ near Y . Since in the decomposition (3.6), the fiber inner product on ΛX -valued forms is $H = K \oplus \tilde{\beta}^{-2}K$ in view of (2.1), we have

$$d_X = \begin{pmatrix} d_Y & 0 \\ \partial_\alpha & -d_Y \end{pmatrix} \quad \text{and} \quad \delta_X = \begin{pmatrix} \delta_Y & \partial_\alpha^* \\ 0 & -\delta_Y \end{pmatrix}, \quad (3.11)$$

where d_Y is the differential on Y and ∂_α^* is the formal adjoint of $\partial_\alpha: \mathcal{C}^\infty(X; \Lambda Y) \subset L^2(X, |dh|; \Lambda Y, K) \rightarrow L^2(X, |dh|; \Lambda Y, \tilde{\beta}^{-2}K)$. Thus, if $\tilde{\beta}$ and k are independent of α near Y , we simply have

$$\partial_\alpha^* = -\beta^{-2}\partial_\alpha,$$

and in general, $\partial_\alpha^* = -\beta^{-2}\partial_\alpha + \alpha^2 p_1 \partial_\alpha + \alpha p_2$, where $p_1, p_2 \in \mathcal{C}^\infty(\overline{X}_{\text{even}})$.

Finally, we compute the form of $\hat{d}(\sigma)$ near Y acting on forms as in (3.10):

$$\hat{d}(\sigma)\mathcal{C} = \begin{pmatrix} d_Y & 0 & 0 & 0 \\ \partial_\alpha & -\alpha d_Y & -\beta\alpha^{-1}d_Y & 0 \\ -i\sigma\alpha^{-1} & 0 & -\alpha^{-1}d_Y & 0 \\ 0 & -i\sigma & -i\sigma\beta\alpha^{-2} - \alpha^{-1}\partial_\alpha & d_Y \end{pmatrix}. \quad (3.12)$$

Thus, applying $\hat{d}(\sigma)$ to $u \in \mathcal{C}_{(\sigma)}^\infty$ yields an element

$$\hat{d}(\sigma)u \in \begin{pmatrix} \alpha^{-i\beta\sigma}\mathcal{C}^\infty(\overline{X}_{\text{even}}; \Lambda Y) \\ \alpha^{-i\beta\sigma-1}\mathcal{C}^\infty(\overline{X}_{\text{even}}; \Lambda Y) \\ \alpha^{-i\beta\sigma-1}\mathcal{C}^\infty(\overline{X}_{\text{even}}; \Lambda Y) \\ \alpha^{-i\beta\sigma}\mathcal{C}^\infty(\overline{X}_{\text{even}}; \Lambda Y) \end{pmatrix},$$

where we use that there is a cancellation in the $(4, 3)$ entry of $\widehat{d}(\sigma)\mathcal{C}$ in view of $(i\sigma\beta\alpha^{-2} + \alpha^{-1}\partial_\alpha)\alpha^{-i\beta\sigma} = 0$; without this cancellation, the fourth component of $\widehat{d}(\sigma)u$ would only lie in $\alpha^{-i\beta\sigma-2}\mathcal{C}^\infty(\overline{X}_{\text{even}}; \Lambda Y)$. Similarly, we compute

$$\widehat{\delta}(\sigma)\mathcal{C} = \begin{pmatrix} -\delta_Y & -\alpha^{-1}\partial_\alpha^*\alpha^2 & -\beta\alpha^{-1}\partial_\alpha^* + i\sigma\alpha^{-2} & 0 \\ 0 & \alpha\delta_Y & \beta\alpha^{-1}\delta_Y & i\sigma\alpha^{-1} \\ 0 & 0 & \alpha^{-1}\delta_Y & \partial_\alpha^* \\ 0 & 0 & 0 & -\delta_Y \end{pmatrix}, \quad (3.13)$$

thus applying $\widehat{\delta}(\sigma)$ to $u \in \mathcal{C}_{(\sigma)}^\infty$ also gives an element

$$\widehat{\delta}(\sigma)u \in \begin{pmatrix} \alpha^{-i\beta\sigma}\mathcal{C}^\infty(\overline{X}_{\text{even}}; \Lambda Y) \\ \alpha^{-i\beta\sigma-1}\mathcal{C}^\infty(\overline{X}_{\text{even}}; \Lambda Y) \\ \alpha^{-i\beta\sigma-1}\mathcal{C}^\infty(\overline{X}_{\text{even}}; \Lambda Y) \\ \alpha^{-i\beta\sigma}\mathcal{C}^\infty(\overline{X}_{\text{even}}; \Lambda Y) \end{pmatrix},$$

where there is again a cancellation in the $(1, 3)$ entry of $\widehat{\delta}(\sigma)\mathcal{C}$; without this cancellation, the first component of $\widehat{d}(\sigma)u$ would only lie in $\alpha^{-i\beta\sigma-2}\mathcal{C}^\infty(\overline{X}_{\text{even}}; \Lambda Y)$.

In fact, a bit more is true: Namely, one checks that⁶ $\alpha^{i\beta\sigma}\mathcal{C}^{-1}\widehat{d}(\sigma)\mathcal{C}\alpha^{-i\beta\sigma}$ and $\alpha^{i\beta\sigma}\mathcal{C}^{-1}\widehat{\delta}(\sigma)\mathcal{C}\alpha^{-i\beta\sigma}$ preserve the space $\mathcal{C}^\infty(\overline{X}_{\text{even}}; \Lambda Y)^4$ (in the decomposition (3.8)), hence if $u \in \mathcal{C}_{(\sigma)}^\infty$, then also $\widehat{d}(\sigma)u, \widehat{\delta}(\sigma)u \in \mathcal{C}_{(\sigma)}^\infty$. Since it will be useful later, we check this explicitly for $\sigma = 0$ by computing

$$\mathcal{C}^{-1}\widehat{d}(0)\mathcal{C} = \begin{pmatrix} d_Y & 0 & 0 & 0 \\ \alpha^{-1}\partial_\alpha & -d_Y & 0 & 0 \\ 0 & 0 & -d_Y & 0 \\ 0 & 0 & -\alpha^{-1}\partial_\alpha & d_Y \end{pmatrix} \quad (3.14)$$

and

$$\mathcal{C}^{-1}\widehat{\delta}(0)\mathcal{C} = \begin{pmatrix} -\delta_Y & -\alpha^{-1}\partial_\alpha^*\alpha^2 & -\alpha^{-1}\partial_\alpha^*\beta & 0 \\ 0 & \delta_Y & 0 & -\beta\alpha^{-1}\partial_\alpha^* \\ 0 & 0 & \delta_Y & \alpha\partial_\alpha^* \\ 0 & 0 & 0 & -\delta_Y \end{pmatrix}. \quad (3.15)$$

3.1. Absence of resonances in $\text{Im } \sigma > 0$. The fiber inner product on the form bundle is not positive definite, thus we cannot use standard arguments for (formally) self-adjoint operators to exclude a non-trivial kernel of $\widehat{d}(\sigma) + \widehat{\delta}(\sigma)$. We therefore introduce a different inner product (by which we mean here a non-degenerate sesquilinear form), related to the natural inner product induced by the metric, which does have some positivity properties. Concretely, for $\theta \in (-\pi/2, \pi/2)$, we use the inner product $H \oplus e^{-2i\theta}H$, i.e. on pure degree k -forms on M , the fiber inner product is given by $H_k \oplus e^{-2i\theta}H_{k-1}$ in the decomposition into tangential and normal components as in (3.1).

Lemma 3.3. *Let $\theta \in (-\pi/2, \pi/2)$. Suppose that $u \in L^2(\alpha|dh|; H \oplus H)$ is such that $\langle u, u \rangle_{H \oplus e^{-2i\theta}H} = 0$. Then $u = 0$.*

⁶Either, this follows by a direct computation; or one notes that these operators are equal (up to a smooth phase factor) to the matrices of the Fourier transforms in t_* of d and δ with respect to the form decomposition (3.8), which are smooth on the extended manifold \tilde{X} .

Proof. With $u = u_T + \alpha dt \wedge u_N$, we have $\|u_T\|_{L^2(\alpha|dh|;H)}^2 + e^{-2i\theta}\|u_N\|_{L^2(\alpha|dh|;H)}^2 = 0$. Multiplying this equation by $e^{i\theta}$ and taking real parts gives

$$\cos(\theta)\|u\|_{L^2(\alpha|dh|;H \oplus H)}^2 = 0,$$

hence $u = 0$, since $\cos \theta > 0$ for θ in the given range. \square

Using the volume density $\alpha|dh|$ to compute adjoints,⁷ we have

$$\langle \widehat{d}(\sigma)u, v \rangle_{H \oplus e^{-2i\theta}H} = \langle u, \widehat{\delta}_\theta(\sigma)v \rangle_{H \oplus e^{-2i\theta}H}, \quad u, v \in \mathcal{C}_c^\infty(X; \Lambda X \oplus \Lambda X)$$

for the operator

$$\widehat{\delta}_\theta(\sigma) = \begin{pmatrix} \alpha^{-1}\delta_X\alpha & ie^{2i\theta}\bar{\sigma}\alpha^{-1} \\ 0 & -\delta_X \end{pmatrix},$$

which equals $-\widehat{\delta}(\sigma)$ provided $e^{2i\theta}\bar{\sigma} = -\sigma$, i.e. $\sigma \in e^{i\theta} \cdot i(0, \infty)$.

Remark 3.4. Since the inner product $H \oplus e^{-2i\theta}H$ is not Hermitian, we do *not* have $\langle \widehat{\delta}_\theta(\sigma)u, v \rangle_{H \oplus e^{-2i\theta}H} = \langle u, \widehat{d}(\sigma)v \rangle_{H \oplus e^{-2i\theta}H}$ in general. Rather, one computes

$$\begin{aligned} \langle \widehat{\delta}_\theta(\sigma)u, v \rangle_{H \oplus e^{2i\theta}H} &= \overline{\langle v, \widehat{\delta}_\theta(\sigma)u \rangle_{H \oplus e^{-2i\theta}H}} \\ &= \overline{\langle \widehat{d}(\sigma)v, u \rangle_{H \oplus e^{-2i\theta}H}} = \langle u, \widehat{d}(\sigma)v \rangle_{H \oplus e^{2i\theta}H}. \end{aligned} \quad (3.16)$$

Now suppose $u \in \mathcal{C}_{(\sigma)}^\infty$ is a solution, with $\text{Im } \sigma > 0$, of

$$(\widehat{d}(\sigma) + \widehat{\delta}(\sigma))u = 0. \quad (3.17)$$

We claim that every such u must vanish. To show this, we apply $\widehat{d}(\sigma)$ to (3.17) and pair the result with u ; this gives

$$\begin{aligned} 0 &= \langle \widehat{d}(\sigma)\widehat{\delta}(\sigma)u, u \rangle_{H \oplus e^{-2i\theta}H} = \langle \widehat{\delta}(\sigma)u, \widehat{\delta}_\theta(\sigma)u \rangle_{H \oplus e^{-2i\theta}H} \\ &= -\langle \widehat{\delta}(\sigma)u, \widehat{d}(\sigma)u \rangle_{H \oplus e^{-2i\theta}H}, \end{aligned} \quad (3.18)$$

where we choose $\theta \in (-\pi/2, \pi/2)$ so that $\sigma \in e^{i\theta} \cdot i(0, \infty)$; the integration by parts will be justified momentarily. By Lemma 3.3, this implies $\widehat{\delta}(\sigma)u = 0$. On the other hand, applying $\widehat{\delta}(\sigma)$ to (3.17) and using (3.16), we get, for $\sigma \in e^{i\theta} \cdot i(0, \infty)$,

$$\begin{aligned} 0 &= \langle \widehat{\delta}(\sigma)\widehat{d}(\sigma)u, u \rangle_{H \oplus e^{2i\theta}H} = -\langle \widehat{\delta}_\theta(\sigma)\widehat{d}(\sigma)u, u \rangle_{H \oplus e^{2i\theta}H} \\ &= -\langle \widehat{d}(\sigma)u, \widehat{d}(\sigma)u \rangle_{H \oplus e^{2i\theta}H}, \end{aligned} \quad (3.19)$$

hence $\widehat{d}(\sigma)u = 0$ by Lemma 3.3, again modulo justifying the integration by parts.

Using the splitting (3.1) and the form (3.5) of $\widehat{d}(\sigma)$, the second component of the equation $\widehat{d}(\sigma)u = 0$ gives $i\sigma u_T + d_X \alpha u_N = 0$. Taking the $L^2(\alpha|dh|; H)$ -pairing of this with u_T gives (the integration by parts to be justified below)

$$0 = i\sigma\|u_T\|^2 + \langle d_X \alpha u_N, u_T \rangle = i\sigma\|u_T\|^2 + \langle u_N, \delta_X \alpha u_T \rangle, \quad (3.20)$$

and then the first component of $\widehat{\delta}(\sigma)u = 0$, i.e. $\delta_X \alpha u_T = i\sigma u_N$, can be used to rewrite the pairing on the right hand side; we obtain $0 = i(\sigma\|u_T\|^2 - \bar{\sigma}\|u_N\|^2)$. Writing $\sigma = ie^{i\theta}\tilde{\sigma}$ with $\tilde{\sigma} > 0$ real, this becomes

$$0 = \tilde{\sigma}(e^{i\theta}\|u_T\|^2 + e^{-i\theta}\|u_N\|^2), \quad (3.21)$$

and taking the real part of this equation gives $u_T = 0 = u_N$, hence $u = 0$.

⁷See Definition 3.1 for the notation used here.

We now justify the integrations by parts used in (3.18) and (3.19), which is only an issue at Y . First of all, since $u \in \mathcal{C}_{(\sigma)}^\infty$ and $\text{Im } \sigma > 0$, the pairings are well-defined in the strong sense that all functions which appear in the pairings are elements of $L^2(\alpha|dh|; H \oplus H)$; in fact, all functions in these pairings lie in $\mathcal{C}_{(\sigma)}^\infty$. In view of the block structure $H \oplus e^{-2i\theta}H = K \oplus \tilde{\beta}^{-2}K \oplus e^{-2i\theta}K \oplus \tilde{\beta}^{-2}e^{-2i\theta}K$ of the inner product, the only potentially troublesome term for the integration by parts there is the pairing of the first components, since this is where we need the cancellation mentioned after (3.13) to ensure that $\hat{d}(\sigma)u \in L^2$. However, if we only use the cancellation in one of the terms, we pair $\alpha^{-i\beta\sigma}\mathcal{C}^\infty(\bar{X}_{\text{even}}; \Lambda Y)$ against $\alpha^{-i\beta\sigma-2}\mathcal{C}^\infty(\bar{X}_{\text{even}}; \Lambda Y)$ in the first component, thus this pairing is still absolutely integrable and one can integrate by parts. Likewise, the integration by parts used in (3.19) only has potential issues in the pairing of the fourth components, since we need the cancellation mentioned after (3.12) to ensure that $\hat{d}(\sigma)u \in L^2$. But again, if we only use this cancellation in one of the terms, we pair $\alpha^{-i\beta\sigma}\mathcal{C}^\infty(\bar{X}_{\text{even}}; \Lambda Y)$ against $\alpha^{-i\beta\sigma-2}\mathcal{C}^\infty(\bar{X}_{\text{even}}; \Lambda Y)$, which is absolutely integrable.

In order to justify (3.20), we observe using (3.11) that near Y ,

$$u_T, d_X \alpha u_N \in \begin{pmatrix} \alpha^{-i\beta\sigma}\mathcal{C}^\infty \\ \alpha^{-i\beta\sigma-1}\mathcal{C}^\infty \end{pmatrix}, \quad u_N, \delta_X \alpha u_T \in \begin{pmatrix} \alpha^{-i\beta\sigma-1}\mathcal{C}^\infty \\ \alpha^{-i\beta\sigma}\mathcal{C}^\infty \end{pmatrix},$$

where we write $\mathcal{C}^\infty = \mathcal{C}^\infty(\bar{X}_{\text{even}}; \Lambda Y)$. These membership statements do not rely on any cancellations, and since all these functions are in $L^2(\alpha|dh|; \Lambda Y, K)$ near Y , the integration by parts in (3.20) is justified.

We summarize the above discussion and extend it to a quantitative version:

Proposition 3.5. *There exists a constant $C > 0$ such that for all $\sigma \in \mathbb{C}$ with $\text{Im } \sigma > 0$, we have the following estimate for $u \in \mathcal{C}_{(\sigma)}^\infty$:*

$$\|u\|_{L^2(\alpha|dh|; H \oplus H)} \leq C \frac{|\sigma|}{|\text{Im } \sigma|^2} \|(\hat{d}(\sigma) + \hat{\delta}(\sigma))u\|_{L^2(\alpha|dh|; H \oplus H)}. \quad (3.22)$$

Proof. Write $\sigma = ie^{i\theta}\tilde{\sigma}$, $\theta \in (-\pi/2, \pi/2)$, $\tilde{\sigma} > 0$, as before. Let $f = (\hat{d}(\sigma) + \hat{\delta}(\sigma))u$; in particular $f \in \mathcal{C}_{(\sigma)}^\infty$. Then $\hat{d}(\sigma)\hat{\delta}(\sigma)u = \hat{d}(\sigma)f$, so

$$\langle \hat{\delta}(\sigma)u, \hat{\delta}(\sigma)u \rangle_{H \oplus e^{-2i\theta}H} = -\langle \hat{d}(\sigma)\hat{\delta}(\sigma)u, u \rangle_{H \oplus e^{-2i\theta}H} = \langle f, \hat{\delta}(\sigma)u \rangle_{H \oplus e^{-2i\theta}H}, \quad (3.23)$$

and similarly

$$\langle \hat{d}(\sigma)u, \hat{d}(\sigma)u \rangle_{H \oplus e^{2i\theta}H} = \langle f, \hat{d}(\sigma)u \rangle_{H \oplus e^{2i\theta}H}. \quad (3.24)$$

Multiply (3.23) by $e^{i\theta}$, (3.24) by $e^{-i\theta}$ and take the sum of both equations to get

$$\begin{aligned} e^{i\theta}(\|\hat{\delta}(\sigma)u\|_T^2 + \|\hat{d}(\sigma)u\|_N^2) + e^{-i\theta}(\|\hat{\delta}(\sigma)u\|_N^2 + \|\hat{d}(\sigma)u\|_T^2) \\ = e^{i\theta}\langle f, \hat{\delta}(\sigma)u \rangle_{H \oplus e^{-2i\theta}H} + e^{-i\theta}\langle f, \hat{d}(\sigma)u \rangle_{H \oplus e^{2i\theta}H}. \end{aligned}$$

Here, the norms without subscript are $L^2(\alpha|dh|; H \oplus H)$ -norms as usual. Taking the real part and applying Cauchy-Schwartz to the right hand side produces the estimate

$$\|\hat{d}(\sigma)u\| + \|\hat{\delta}(\sigma)u\| \leq \frac{4}{\cos \theta} \|f\| = \frac{4|\sigma|}{|\text{Im } \sigma}| \|f\|. \quad (3.25)$$

We estimate u in terms of the left hand side of (3.25) by following the arguments leading to (3.21): Put $v = \hat{d}(\sigma)u$ and $w = \hat{\delta}(\sigma)u$. Then $i\sigma u_T + d_X \alpha u_N = -\alpha v_N$;

we pair this with u_T in $L^2(\alpha|dh|; H)$ and obtain

$$i\sigma\|u_T\|^2 + \langle u_N, \delta_X \alpha u_T \rangle = -\langle \alpha v_N, u_T \rangle.$$

Using $-\delta_X \alpha u_T + i\sigma u_N = \alpha w_T$, this implies

$$i\sigma\|u_T\|^2 - i\bar{\sigma}\|u_N\|^2 = -\langle \alpha v_N, u_T \rangle + \langle u_N, \alpha w_T \rangle,$$

thus

$$\tilde{\sigma}(e^{i\theta}\|u_T\|^2 + e^{-i\theta}\|u_N\|^2) = \langle \alpha v_N, u_T \rangle - \langle u_N, \alpha w_T \rangle.$$

Taking the real part and applying Cauchy-Schwartz, we get

$$(\cos \theta)\|u\| \leq |\sigma|^{-1}(\|\alpha v\| + \|\alpha w\|) \lesssim |\sigma|^{-1}(\|v\| + \|w\|).$$

In combination with (3.25), this yields (3.22). \square

3.2. Boundary pairing and absence of non-zero real resonances. We proceed to exclude non-zero real resonances for $d + \delta$ by means of a boundary pairing argument similar to [41, §2.3].

Proposition 3.6. *Suppose $\sigma \in \mathbb{R}$, $\sigma \neq 0$. If $u \in \mathcal{C}_{(\sigma)}^\infty$ solves $(\widehat{d}(\sigma) + \widehat{\delta}(\sigma))u = 0$, then $u = 0$.*

Proof. Writing $u = u_T + \alpha dt \wedge u_N$ as usual, we can expand $(\widehat{d}(\sigma) + \widehat{\delta}(\sigma))u = 0$ as

$$\begin{aligned} (\alpha d_X - \delta_X \alpha)u_T + i\sigma u_N &= 0 \\ -i\sigma u_T + (-d_X \alpha + \alpha \delta_X)u_N &= 0. \end{aligned} \quad (3.26)$$

Applying $(-d_X \alpha + \alpha \delta_X)$ to the first equation and using the second equation to simplify the resulting expression produces a second order equation for u_T ,

$$(d_X \alpha \delta_X \alpha + \alpha \delta_X \alpha d_X - d_X \alpha^2 d_X - \sigma^2)u_T = 0. \quad (3.27)$$

Writing $u_T = u_{TT} + d\alpha \wedge u_{TN}$ as in (3.6), we see from the definition of the space $\mathcal{C}_{(\sigma)}^\infty$ that

$$u_T \in \mathcal{C}_{(\sigma),T}^\infty := \alpha^{-i\beta\sigma} \mathcal{C}^\infty(\overline{X}_{\text{even}}; \Lambda Y) \oplus \alpha^{-i\beta\sigma-1} \mathcal{C}^\infty(\overline{X}_{\text{even}}; \Lambda Y)$$

near Y . Notice that the space $\mathcal{C}_{(\sigma),T}^\infty$ barely fails to be contained in $L^2(\alpha|dh|)$.

We will deduce from (3.27) that $u_T = 0$; equation (3.26) then gives $u_N = 0$, as $\sigma \neq 0$. Now, the $L^2(\alpha|dh|; H)$ -adjoint of $d_X \alpha$ is $\delta_X \alpha$, hence even ignoring the term $d_X \alpha^2 d_X$, the operator in (3.27) is not symmetric. However, we can obtain a simpler equation from (3.27) by applying d_X to it; write $v_T = d_X u_T \in \mathcal{C}_{(\sigma),T}^\infty$, and near Y ,

$$v_T = \begin{pmatrix} \alpha^{-i\beta\sigma} \tilde{v}_{TT} \\ \alpha^{-i\beta\sigma-1} \tilde{v}_{TN} \end{pmatrix}, \quad \tilde{v}_{TT}, \tilde{v}_{TN} \in \mathcal{C}^\infty(\overline{X}_{\text{even}}; \Lambda Y).$$

Then v_T satisfies the equation

$$(d_X \alpha \delta_X \alpha - \sigma^2)v_T = 0,$$

and $d_X \alpha \delta_X \alpha$ is symmetric with respect to the $L^2(\alpha|dh|; H)$ -inner product. We now compute the boundary pairing formula (using the same inner product); to this end, pick a cutoff function $\chi \in \mathcal{C}^\infty(\overline{X})$ such that in a collar neighborhood $[0, \delta)_\alpha \times Y_y$ of Y in \overline{X} , $\chi = \chi(\alpha)$ is identically 0 near $\alpha = 0$ and identically 1 in $\alpha \geq \delta/2$, and extend χ by 1 to all of \overline{X} . Define $\chi_\epsilon(\alpha) = \chi(\alpha/\epsilon)$ and $\chi'_\epsilon(\alpha) = \chi'(\alpha/\epsilon)$. Then

$$\begin{aligned} 0 &= \lim_{\epsilon \rightarrow 0} (\langle (d_X \alpha \delta_X \alpha - \sigma^2)v_T, \chi_\epsilon v_T \rangle - \langle v_T, \chi_\epsilon (d_X \alpha \delta_X \alpha - \sigma^2)v_T \rangle) \\ &= \lim_{\epsilon \rightarrow 0} \langle v_T, [d_X \alpha \delta_X \alpha, \chi_\epsilon] v_T \rangle. \end{aligned} \quad (3.28)$$

The coefficients of the commutator are supported near Y , hence we use (3.11) to compute its form as

$$\begin{aligned} [d_X \alpha \delta_X \alpha, \chi_\epsilon] &= \left[\begin{pmatrix} d_Y \alpha \delta_Y \alpha & d_Y \alpha \partial_\alpha^* \alpha \\ \partial_\alpha \alpha \delta_Y \alpha & \partial_\alpha \alpha \partial_\alpha^* \alpha + d_Y \alpha \delta_Y \alpha \end{pmatrix}, \chi_\epsilon \right] \\ &= \begin{pmatrix} 0 & d_Y \alpha [\partial_\alpha^*, \chi_\epsilon] \alpha \\ [\partial_\alpha, \chi_\epsilon] \alpha \delta_Y \alpha & [\partial_\alpha \alpha \partial_\alpha^* \alpha, \chi_\epsilon] \end{pmatrix} \\ &= \epsilon^{-1} \begin{pmatrix} 0 & -\beta^{-2}(\alpha^2 + \mathcal{O}(\alpha^4)) \chi'_\epsilon d_Y \\ \chi'_\epsilon \alpha \delta_Y \alpha & \chi'_\epsilon \alpha \partial_\alpha^* \alpha - \partial_\alpha(\alpha^2 + \mathcal{O}(\alpha^4)) \beta^{-2} \chi'_\epsilon \end{pmatrix}. \end{aligned}$$

In (3.28), the off-diagonal terms of this give terms of the form

$$\int_Y \int \alpha^{\mp i\beta\sigma} \alpha^{\pm i\beta\sigma-1} \epsilon^{-1} \alpha^2 \chi'_\epsilon \tilde{v} d\alpha |dk| \quad (3.29)$$

with $\tilde{v} \in \mathcal{C}^\infty(\overline{X}_{\text{even}})$, and are easily seen to vanish in the limit $\epsilon \rightarrow 0$. The non-zero diagonal term gives⁸

$$\begin{aligned} &\epsilon^{-1} \langle \alpha^{-i\beta\sigma-1} \tilde{v}_{TN}, (\chi'_\epsilon \alpha \partial_\alpha^* \alpha - \partial_\alpha \alpha^2 \beta^{-2} \chi'_\epsilon) \alpha^{-i\beta\sigma-1} \tilde{v}_{TN} \rangle_{L^2(X; \alpha\beta d\alpha |dk|; \Lambda Y; \beta^{-2} K)} \\ &= 2 \int_Y \int \langle \tilde{v}_{TN}, i\beta^{-2} \sigma \tilde{v}_{TN} \rangle_K \epsilon^{-1} \chi'_\epsilon d\alpha |dk| + o(1) \\ &\xrightarrow{\epsilon \rightarrow 0} -2i\beta^{-2} \sigma \|\tilde{v}_{TN}|_Y\|_{L^2(Y, |dk|; K)}^2; \end{aligned}$$

here, both summands in the pairing yield the same result, as is most easily seen by integrating by parts in α , hence the factor of 2, and the $o(1)$ -term comes from differentiating \tilde{v}_{TN} , which produces a term of the form (3.29). We thus arrive at

$$0 = \langle (d_X \alpha \delta_X \alpha - \sigma^2) v_T, v_T \rangle - \langle v_T, (d_X \alpha \delta_X \alpha - \sigma^2) v_T \rangle = -2i\beta^{-2} \sigma \|\tilde{v}_{TN}|_Y\|^2,$$

whence $\tilde{v}_{TN}|_Y = 0$ in view of $\sigma \neq 0$, so we in fact have

$$v_T = \begin{pmatrix} \alpha^{-i\beta\sigma} \tilde{v}_{TT} \\ \alpha^{-i\beta\sigma} \tilde{v}'_{TN} \end{pmatrix}, \quad \tilde{v}'_{TN} \in \mathcal{C}^\infty(\overline{X}_{\text{even}}; \Lambda Y). \quad (3.30)$$

For the next step, recall that on a manifold with boundary \overline{X} , 0-vector fields, introduced by Mazzeo and Melrose [39] to analyze the resolvent of the Laplacian on asymptotically hyperbolic spaces, are smooth vector fields that vanish at ∂X , i.e. are of the form αV , where V is a smooth vector field on \overline{X} , and α , as in our case, is a boundary defining function, i.e. in local coordinates a linear combination, with smooth coefficients, of $\alpha \partial_\alpha$ and $\alpha \partial_{y_j}$. Further, 0-differential operators, $A \in \text{Diff}_0(\overline{X})$, are the differential operators generated by these (taking finite sums of finite products, with $\mathcal{C}^\infty(\overline{X})$ coefficients). As a contrast, b-vector fields are merely tangent to ∂X , so in local coordinates they are linear combinations, with smooth coefficients, of $\alpha \partial_\alpha$ and ∂_{y_j} , and they generate b-differential operators $\text{Diff}_b(\overline{X})$. Often, as in our case, one is considering solutions of 0-differential equations with additional properties, such as having an expansion in powers of α (and perhaps $\log \alpha$) with smooth coefficients, i.e. polyhomogeneous functions. In these cases $\alpha \text{Diff}_b(\overline{X}) \subset \text{Diff}_0(\overline{X})$ acts ‘trivially’ on an expansion in that it maps each term to one with an additional order of vanishing, so in particular, one can analyze the asymptotic expansion of solutions of 0-differential equations in this restrictive class by ignoring the $\alpha \text{Diff}_b(\overline{X})$ terms. Notice that $\alpha \partial_{y_j} \in \alpha \text{Diff}_b(\overline{X})$ in particular, so

⁸Recall that the volume density is given by $\alpha |dh| = \alpha \beta d\alpha |dk|$, and the fiber inner product in the (TN) -component is $\beta^{-2} K$.

the tangential 0-derivatives can be dropped for this purpose. The indicial equation is then obtained by freezing the coefficients of $A \in \text{Diff}_0(\bar{X})$ at ∂X , i.e. writing it as $\sum_{k,\beta} a_{k,\beta}(\alpha, y)(\alpha\partial_\alpha)^k(\alpha\partial_y)^\beta$, where $a_{k,\beta}$ are bundle endomorphism valued, and restricting α to 0, and dropping all terms with a positive power of $\alpha\partial_y$, to obtain $\sum_k a_{k,0}(0, y)(\alpha\partial_\alpha)^k$. This can be thought of as a regular-singular ODE in α for each y ; its indicial roots are called the indicial roots of the original 0-operator, and they determine the asymptotics of solutions of the homogeneous PDE with this a priori form.

Now $d_X\alpha\delta_X\alpha - \sigma^2 \in \text{Diff}_0^2(\bar{X})$ is a 0-differential operator which equals

$$d_X\alpha\delta_X\alpha - \sigma^2 = \begin{pmatrix} -\sigma^2 & 0 \\ 0 & -\beta^{-2}\partial_\alpha\alpha\partial_\alpha\alpha - \sigma^2 \end{pmatrix}$$

modulo $\alpha\text{Diff}_b^2(\bar{X})$; hence its indicial roots are $\pm i\beta\sigma - 1$. In particular, $-i\beta\sigma + j$, $j \in \mathbb{N}_0$, is not an indicial root. Thus, a standard inductive argument starting with (3.30) shows that $v_T \in \dot{\mathcal{C}}^\infty(\bar{X}; \Lambda X)$.

Next, we note that v_T lies in the kernel of the operator

$$d_X\alpha\delta_X\alpha + \alpha^{-1}\delta_X\alpha^3d_X - \sigma^2 \in \text{Diff}_0^2(\bar{X}; {}^0\Lambda\bar{X}),$$

which has the same principal part as $\alpha^2\Delta_X$, hence is principally a 0-Laplacian; thus, we can apply Mazzeo's result on unique continuation at infinity [38] to conclude that the rapidly vanishing v_T must in fact vanish identically.

We thus have proved $d_Xu_T = 0$. Since u_T satisfies (3.27), we deduce that u_T itself satisfies

$$(d_X\alpha\delta_X\alpha - \sigma^2)u_T = 0,$$

thus repeating the above argument shows that this implies $u_T = 0$, hence $u = 0$, and the proof is complete. \square

3.3. Analysis of the zero resonance. We have shown now that the only potential resonance for $d + \delta$ in $\text{Im } \sigma \geq 0$ is $\sigma = 0$, and we proceed to study the zero resonance in detail, in particular giving a cohomological interpretation of it in Section 3.4.

We begin by establishing the order of the pole of $(\tilde{d}(\sigma) + \tilde{\delta}(\sigma))^{-1}$:

Lemma 3.7. $(\tilde{d}(\sigma) + \tilde{\delta}(\sigma))^{-1}$ has a pole of order 1 at $\sigma = 0$.

Proof. Since $\tilde{d}(0) + \tilde{\delta}(0)$ annihilates constant functions (which are indeed elements of $\mathcal{C}_{(0)}^\infty$), $(\tilde{d}(\sigma) + \tilde{\delta}(\sigma))^{-1}$ does have a pole at 0. Denote the order of the pole by N . Then there is a holomorphic family $\tilde{u}(\sigma) \in \mathcal{C}^\infty(\tilde{X})$ with $\tilde{u}(0) \neq 0$ such that $(\tilde{d}(\sigma) + \tilde{\delta}(\sigma))\tilde{u}(\sigma) = \sigma^N\tilde{v}$, where $\tilde{v} \in \mathcal{C}^\infty(\tilde{X})$. Define $u(\sigma) = e^{iF\sigma}\tilde{u}(\sigma)|_X \in \mathcal{C}_{(\sigma)}^\infty$ and $v(\sigma) = e^{iF\sigma}\tilde{v}|_X \in \mathcal{C}_{(\sigma)}^\infty$, then $(\hat{d}(0) + \hat{\delta}(0))u(\sigma) = \sigma^Nv(\sigma)$. Moreover, since $(\tilde{d}(0) + \tilde{\delta}(0))\tilde{u}(0) = 0$ and $\tilde{u}(0)$ is non-zero, Lemma 2.2 shows that $u(0) \neq 0$.

Let us assume now that $N \geq 2$. For $\sigma \in i(0, \infty)$ close to 0, the quantitative estimate in Proposition 3.5 now gives

$$\|u(\sigma)\| \lesssim |\sigma|^{-1+N}\|v(\sigma)\| \leq |\sigma|\|v(\sigma)\|, \quad (3.31)$$

where we use the norm of $L^2(\alpha|dh|; H \oplus H)$.⁹ Notice that this does not immediately give $u(0) = 0$ since $v(0) \notin L^2(\alpha|dh|; H \oplus H)$. However, we can quantify

⁹Observe that in the notation of Section 3.1, we have $\hat{\delta}_0(0) = -\hat{\delta}(0)$, hence using the Riemannian fiber inner product $H \oplus H$ is natural when studying the zero resonance.

the degeneration of the L^2 -norm of $v(\sigma)$ as $\sigma \rightarrow 0$. To see this, we first observe that the L^2 -norm of $v(\sigma)$ restricted to the complement of any fixed neighborhood of Y does stay bounded, so it remains to analyze the L^2 -norms of the four components of $v(\sigma)$ near Y in the notation of (3.7); denote these components by $\alpha^{-i\beta\sigma}\tilde{v}_{TT}(\sigma)$, $\alpha^{-i\beta\sigma-1}\tilde{v}_{TN}(\sigma)$, $\alpha^{-i\beta\sigma-1}\tilde{v}_{NT}(\sigma)$ and $\alpha^{-i\beta\sigma}\tilde{v}_{NN}(\sigma)$, so that the $\tilde{v}_{\bullet\bullet}(\sigma) \in \mathcal{C}^\infty(\overline{X}_{\text{even}}; \Lambda Y)$ uniformly. Since the fiber metric in this basis has a block diagonal form and any $\mathcal{C}^\infty(\overline{X}_{\text{even}})$ -multiple of $\alpha^{-i\beta\sigma}$ is uniformly square-integrable with respect to the volume density $\alpha|dh|$, the degeneration of the L^2 -norm of v is caused by the (TN) and (NT) components. For these, we compute, with $\tilde{w}(\sigma) \in \mathcal{C}^\infty(\overline{X}_{\text{even}}; \Lambda Y)$ denoting any continuous family supported near Y ,

$$\begin{aligned} & \int_Y \int \alpha^{2(-i\beta\sigma-1)} \|\tilde{w}\|_K^2 \alpha d\alpha |dk| \\ &= \|\tilde{w}(0)\|_{L^2(Y, |dk|; K)}^2 \int \alpha^{-2i\beta\sigma-1} \chi(\alpha) d\alpha + \mathcal{O}(1), \end{aligned}$$

where $\chi \in \mathcal{C}^\infty(\overline{X})$ is a cutoff, equal to 1 near $\alpha = 0$. We can rewrite the integral using an integration by parts, which yields

$$\int \alpha^{-2i\beta\sigma-1} \chi(\alpha) d\alpha = \frac{1}{2i\beta\sigma} \int \alpha^{-2i\beta\sigma} \chi'(\alpha) d\alpha = \mathcal{O}(|\sigma|^{-1}).$$

Therefore, we obtain the bound $\|v(\sigma)\| = \mathcal{O}(|\sigma|^{-1/2})$. Plugging this into (3.31), we conclude using Fatou's Lemma that $u(0) = 0$, which contradicts our assumption that $u(0) \neq 0$. Hence, the order of the pole is $N \leq 1$, but since it is at least 1, it must be equal to 1. \square

Next, we identify the resonant states. *For brevity, we will write $\hat{d} = \hat{d}(0)$, $\hat{\delta} = \hat{\delta}(0)$ and $\hat{\square} = \hat{\square}_g(0)$.*

Proposition 3.8. *$\ker_{\mathcal{C}_{(0)}^\infty}(\hat{d} + \hat{\delta})$ is equal to the space*

$$\mathcal{H} = \{u \in \mathcal{C}_{(0)}^\infty : \hat{d}u = 0, \hat{\delta}u = 0\}. \quad (3.32)$$

Proof. Given $u \in \mathcal{C}_{(0)}^\infty$ with $(\hat{d} + \hat{\delta})u = 0$, we conclude that $\hat{\square}u = 0$, and since $\hat{\square}$ is symmetric on $L^2(\alpha|dh|; H \oplus H)$, we can obtain information about u by a boundary pairing type argument: Concretely, for a cutoff $\chi \in \mathcal{C}^\infty(\overline{X})$ as in the proof of Proposition 3.6, identically 0 near Y , identically 1 outside a neighborhood of Y and a function of α in a collar neighborhood of Y , and with $\chi_\epsilon(\alpha) = \chi(\alpha/\epsilon)$, $\chi'_\epsilon(\alpha) = \chi'(\alpha/\epsilon)$, we have

$$\begin{aligned} 0 &= -\lim_{\epsilon \rightarrow 0} \langle \chi_\epsilon(\hat{d}\hat{\delta} + \hat{\delta}\hat{d})u, u \rangle = \lim_{\epsilon \rightarrow 0} (\langle \hat{\delta}u, \hat{\delta}\chi_\epsilon u \rangle + \langle \hat{d}u, \hat{d}\chi_\epsilon u \rangle) \\ &= \lim_{\epsilon \rightarrow 0} (\|\chi_\epsilon^{1/2}\hat{\delta}u\|^2 + \|\chi_\epsilon^{1/2}\hat{d}u\|^2) + \lim_{\epsilon \rightarrow 0} (\langle \hat{\delta}u, [\hat{\delta}, \chi_\epsilon]u \rangle + \langle \hat{d}u, [\hat{d}, \chi_\epsilon]u \rangle). \end{aligned} \quad (3.33)$$

Since the commutators are supported near Y , we can compute them in the basis (3.7). Let us write $u = \mathcal{C}\tilde{u}$ as in (3.10) with $\sigma = 0$, then in view of (3.13), we have

$$[\hat{\delta}\mathcal{C}, \chi_\epsilon] = \epsilon^{-1}\chi'_\epsilon \begin{pmatrix} 0 & \beta^{-2}\alpha + \mathcal{O}(\alpha^3) & \beta^{-1}\alpha^{-1} + \mathcal{O}(\alpha) & 0 \\ 0 & 0 & 0 & 0 \\ 0 & 0 & 0 & -\beta^{-2} + \mathcal{O}(\alpha^2) \\ 0 & 0 & 0 & 0 \end{pmatrix}, \quad (3.34)$$

and since therefore only the (TT) and (NT) components of $[\widehat{\delta}\mathcal{C}, \chi_\epsilon]\tilde{u}$ are non-zero, we merely compute

$$\begin{aligned} (\widehat{\delta}\mathcal{C}\tilde{u})_{TT} &= -\delta_Y\tilde{u}_{TT} - \alpha^{-1}\partial_\alpha^*\alpha^2\tilde{u}_{TN} - \beta\alpha^{-1}\partial_\alpha^*\tilde{u}_{NT} \\ &\in -\delta_Y\tilde{u}_{TT} + 2\beta^{-2}\tilde{u}_{TN} - \beta\alpha^{-1}\partial_\alpha^*\tilde{u}_{NT} + \alpha\mathcal{C}^\infty(\overline{X}_{\text{even}}; \Lambda Y), \\ (\widehat{\delta}\mathcal{C}\tilde{u})_{NT} &= \alpha^{-1}\delta_Y\tilde{u}_{NT} + \partial_\alpha^*\tilde{u}_{NN} \in \alpha^{-1}\delta_Y\tilde{u}_{NT} + \mathcal{C}^\infty(\overline{X}_{\text{even}}; \Lambda Y). \end{aligned}$$

Notice here that $\alpha^{-1}\partial_\alpha = 2\partial_\mu$ indeed preserves elements of $\mathcal{C}^\infty(\overline{X}_{\text{even}}; \Lambda Y)$. Now in (3.33), the pairing corresponding to the $(1, 2)$ -component of (3.34) is of the form (3.29) (recall that the volume density is $\alpha|dh| = \alpha\beta d\alpha|dk|$) and hence vanishes in the limit $\epsilon \rightarrow 0$, and we conclude that

$$\begin{aligned} \lim_{\epsilon \rightarrow 0} \langle \widehat{\delta}u, [\widehat{\delta}, \chi_\epsilon]u \rangle &= -\langle \delta_Y\tilde{u}_{TT}|_Y, \tilde{u}_{NT}|_Y \rangle + 2\beta^{-2}\langle \tilde{u}_{TN}|_Y, \tilde{u}_{NT}|_Y \rangle \\ &\quad - \beta\langle (\alpha^{-1}\partial_\alpha^*\tilde{u}_{NT})|_Y, \tilde{u}_{NT}|_Y \rangle - \beta^{-1}\langle \delta_Y\tilde{u}_{NT}|_Y, \tilde{u}_{NN}|_Y \rangle, \end{aligned} \quad (3.35)$$

where we use the $L^2(Y, |dk|; K)$ inner product on the right hand side; we absorbed the factor of β from the volume density $\alpha\beta d\alpha|dk|$ into the functions in the pairings.

In a similar vein, we can use (3.12) to compute

$$[\widehat{d}\mathcal{C}, \chi_\epsilon] = \epsilon^{-1}\chi'_\epsilon \begin{pmatrix} 0 & 0 & 0 & 0 \\ 1 & 0 & 0 & 0 \\ 0 & 0 & 0 & 0 \\ 0 & 0 & -\alpha^{-1} & 0 \end{pmatrix} \quad (3.36)$$

and

$$\begin{aligned} (\widehat{d}\mathcal{C}\tilde{u})_{TN} &= \partial_\alpha\tilde{u}_{TT} - \alpha d_Y\tilde{u}_{TN} - \alpha^{-1}d_Y\beta\tilde{u}_{NT} \\ &\in -\beta\alpha^{-1}d_Y\tilde{u}_{NT} + \mathcal{C}^\infty(\overline{X}_{\text{even}}; \Lambda Y), \\ (\widehat{d}\mathcal{C}\tilde{u})_{NN} &= -\alpha^{-1}\partial_\alpha\tilde{u}_{NT} + d_Y\tilde{u}_{NN}. \end{aligned}$$

Correspondingly,

$$\begin{aligned} \lim_{\epsilon \rightarrow 0} \langle \widehat{d}u, [\widehat{d}, \chi_\epsilon]u \rangle &= -\langle d_Y\tilde{u}_{NT}|_Y, \tilde{u}_{TT}|_Y \rangle + \beta^{-1}\langle (\alpha^{-1}\partial_\alpha\tilde{u}_{NT})|_Y, \tilde{u}_{NT}|_Y \rangle \\ &\quad - \beta^{-1}\langle d_Y\tilde{u}_{NN}|_Y, \tilde{u}_{NT}|_Y \rangle, \end{aligned} \quad (3.37)$$

where we again use the $L^2(Y, |dk|; K)$ inner product on the right hand side; notice with regard to the powers of β that on the (TN) and (NN) components, the fiber inner product is $\beta^{-2}K$.

As a consequence of these computations, we conclude that the pairings in (3.33) stay bounded as $\epsilon \rightarrow 0$, hence $\widehat{d}u, \widehat{\delta}u \in L^2(\alpha|dh|; H \oplus H)$ by Fatou's Lemma. Looking at the most singular terms of $\widehat{d}\mathcal{C}\tilde{u}$ and $\widehat{\delta}\mathcal{C}\tilde{u}$ (again using (3.12) and (3.13)), this necessitates

$$d_Y\tilde{u}_{NT}|_Y = 0, \quad \delta_Y\tilde{u}_{NT}|_Y = 0. \quad (3.38)$$

Therefore, taking (3.35) and (3.37) into account, (3.33) simplifies to

$$\begin{aligned} 0 &= \|\widehat{d}u\|^2 + \|\widehat{\delta}u\|^2 + \beta^{-1}\langle (\alpha^{-1}\partial_\alpha\tilde{u}_{NT})|_Y, \tilde{u}_{NT}|_Y \rangle \\ &\quad - \beta\langle (\alpha^{-1}\partial_\alpha^*\tilde{u}_{NT})|_Y, \tilde{u}_{NT}|_Y \rangle + 2\beta^{-2}\langle \tilde{u}_{TN}|_Y, \tilde{u}_{NT}|_Y \rangle. \end{aligned} \quad (3.39)$$

Moreover, the fourth component of the equation $(\widehat{d} + \widehat{\delta})\mathcal{C}\tilde{u} = 0$ yields

$$-(\alpha^{-1}\partial_\alpha\tilde{u}_{NT})|_Y + d_Y\tilde{u}_{NN}|_Y - \delta_Y\tilde{u}_{NN}|_Y = 0,$$

which we can pair with $\tilde{u}_{NT}|_Y$ relative to $L^2(Y, |dk|; K)$, and then an integration by parts together with (3.38) shows that the first boundary pairing in (3.39) vanishes. Likewise, the first component of $(\hat{d} + \hat{\delta})\mathcal{C}\tilde{u} = 0$ gives

$$d_Y \tilde{u}_{TT}|_Y - \delta_Y \tilde{u}_{TT}|_Y + 2\beta^{-2} \tilde{u}_{TN}|_Y - \beta(\alpha^{-1} \partial_\alpha^* \tilde{u}_{NT})|_Y = 0,$$

which we can again pair with $\tilde{u}_{NT}|_Y$, and in view of (3.38), we conclude that the second line of (3.39) vanishes as well. Thus, finally, (3.39) implies that $\hat{d}u = 0$ and $\hat{\delta}u = 0$.

Conversely, every $u \in \mathcal{C}_{(0)}^\infty$ satisfying $\hat{d}u = 0$ and $\hat{\delta}u = 0$ trivially lies in the kernel of $\hat{d} + \hat{\delta}$. \square

The above proof in particular shows:

Corollary 3.9. *Let $u = \mathcal{C}\tilde{u} \in \mathcal{C}_{(0)}^\infty$ be such that $\hat{d}\hat{\delta}u = 0$ (resp. $\hat{\delta}\hat{d}u = 0$), and assume that $\tilde{u}_{NT}|_Y = 0$.¹⁰ Then $\hat{\delta}u = 0$ (resp. $\hat{d}u = 0$). In particular, $\ker_{\mathcal{C}_{(0)}^\infty \cap L^2} \hat{\square} = \mathcal{H} \cap L^2$.*

Proof. Suppose $\hat{d}\hat{\delta}u = 0$. With a cutoff function χ_ϵ as above, we obtain

$$0 = -\lim_{\epsilon \rightarrow 0} \langle \chi_\epsilon \hat{d}\hat{\delta}u, u \rangle = \lim_{\epsilon \rightarrow 0} \|\chi_\epsilon^{1/2} \hat{\delta}u\|^2 + \lim_{\epsilon \rightarrow 0} \langle \hat{\delta}u, [\hat{\delta}, \chi_\epsilon]u \rangle.$$

In view of (3.35) and $\tilde{u}_{NT}|_Y = 0$, the second term on the right hand side vanishes, and we deduce $\hat{\delta}u = 0$. The proof that $\hat{\delta}\hat{d}u = 0$ implies $\hat{d}u = 0$ is similar and uses (3.37). \square

Corollary 3.10. *We have $\ker \hat{\square} = \ker \hat{d}\hat{\delta} \cap \ker \hat{\delta}\hat{d}$.*

Proof. If $u \in \ker \hat{\square}$, then $(\hat{d} + \hat{\delta})u \in \mathcal{H}$, thus $\hat{\delta}(\hat{d} + \hat{\delta})u = \hat{\delta}\hat{d}u = 0$ and $\hat{d}\hat{\delta}u = 0$. \square

We record another setting in which the boundary terms in the proof of Proposition 3.8 vanish:

Lemma 3.11. *Suppose $v \in \mathcal{C}_{(0)}^\infty$ is a solution of $\hat{\delta}\hat{d}\hat{\delta}v = 0$. Then $\hat{d}\hat{\delta}v = 0$. Likewise, if $v \in \mathcal{C}_{(0)}^\infty$ is a solution of $\hat{d}\hat{\delta}\hat{d}v = 0$, then $\hat{\delta}\hat{d}v = 0$.*

Proof. Write $w = \hat{\delta}v \in \mathcal{C}_{(0)}^\infty$. Then $\hat{\delta}\hat{d}w = 0$ implies, by the proof of Proposition 3.8, that $\hat{d}w \in L^2(\alpha|dh|; H \oplus H)$. Writing $w = \mathcal{C}\tilde{w}$, this in particular implies $d_Y \tilde{w}_{NT}|_Y = 0$; but writing $v = \mathcal{C}\tilde{v}$, we have

$$\tilde{w}_{NT} = (\mathcal{C}^{-1} \hat{\delta} \mathcal{C} \tilde{v})_{NT} = \delta_Y \tilde{v}_{NT} + \alpha \partial_\alpha^* \tilde{v}_{NN},$$

as follows from (3.15). Restricting to Y , we thus have $\tilde{w}_{NT}|_Y = \delta_Y \tilde{v}_{NT}|_Y$, and hence $0 = d_Y \delta_Y \tilde{v}_{NT}|_Y$. We pair this in $L^2(Y, |dk|; K)$ with \tilde{v}_{NT} and integrate by parts, obtaining $\delta_Y \tilde{v}_{NT}|_Y = 0$. But this implies that $\tilde{w}_{NT}|_Y = 0$. By Corollary 3.9, this gives $\hat{d}w = \hat{d}\hat{\delta}v = 0$.

For the second part, we proceed analogously: Letting $w = \hat{d}v \in \mathcal{C}_{(0)}^\infty$, we have $\hat{d}\hat{\delta}w = 0$, thus $\hat{\delta}w \in L^2$. This gives $\delta_Y \tilde{w}_{NT}|_Y = 0$; but by (3.14), $\tilde{w}_{NT}|_Y = -d_Y \tilde{v}_{NT}|_Y$, therefore $\delta_Y \tilde{w}_{NT}|_Y = 0$ implies $d_Y \tilde{v}_{NT}|_Y = 0$, so $\tilde{w}_{NT}|_Y = 0$, which in turn gives $\hat{\delta}w = 0$ by Corollary 3.9, hence $\hat{\delta}\hat{d}v = 0$. \square

¹⁰The latter is equivalent to assuming $u \in L^2(\alpha|dh|)$.

3.4. Cohomological interpretation of zero resonant states. *In this section, we will always work with $\sigma = 0$ and hence simply write $\widehat{d} = \widehat{d}(0)$, $\widehat{\delta} = \widehat{\delta}(0)$, $\widetilde{d} = \widetilde{d}(0)$, $\widetilde{\delta} = \widetilde{\delta}(0)$, $\widehat{\square} = \widehat{\square}(0)$ and $\widetilde{\square} = \widetilde{\square}(0)$.*

The space \mathcal{H} defined in Proposition 3.8 is graded by the form degree, i.e.

$$\mathcal{H} = \bigoplus_{k=0}^n \mathcal{H}^k, \quad (3.40)$$

where \mathcal{H}^k is the space of all $u \in \mathcal{H}$ of pure form degree k . In the decomposition (3.1), this means that u_T is a differential k -form on X , and u_N is a differential $(k-1)$ -form. Likewise, $\mathcal{K} := \ker \widehat{\square}$ is graded by form degree, and we write

$$\ker_{\mathcal{C}_{(0)}^\infty} \widehat{\square} = \bigoplus_{k=0}^n \mathcal{K}^k. \quad (3.41)$$

We aim to relate the spaces \mathcal{H}^k and \mathcal{K}^k to certain cohomology groups associated with \bar{X} . As in the Riemannian setting, the central tool is a Hodge type decomposition adapted to \widehat{d} and $\widehat{\delta}$:

Lemma 3.12. *The following Hodge type decomposition holds on X :*

$$\mathcal{C}_{(0)}^\infty = \ker_{\mathcal{C}_{(0)}^\infty} \widehat{\square} \oplus \text{ran}_{\mathcal{C}_{(0)}^\infty} \widehat{\square}. \quad (3.42)$$

Proof. We first claim that such a decomposition holds on \widetilde{X} , i.e. we claim that

$$\mathcal{C}^\infty(\widetilde{X}) = \ker \widetilde{\square} \oplus \text{ran } \widetilde{\square}. \quad (3.43)$$

First of all, since $\widetilde{\square}$ is Fredholm with index 0, its range is closed, and the codimension of the range equals the dimension of the kernel. Hence, in order to show (3.43), we merely need to check that the intersection of $\ker \widetilde{\square}$ and $\text{ran } \widetilde{\square}$ is trivial. Thus, let $\widetilde{u} \in \ker \widetilde{\square} \cap \text{ran } \widetilde{\square}$, and write $\widetilde{u} = \widetilde{\square} \widetilde{v}$. Let $v = \widetilde{v}|_X$. Then $\widetilde{u} \in \ker \widetilde{\square}$ means, restricting to X and using Corollary 3.10, that $\widehat{d} \widehat{\delta} \widehat{d} \widehat{\delta} v = 0$ and $\widehat{\delta} \widehat{d} \widehat{\delta} \widehat{d} v = 0$. Repeated application of Lemma 3.11 thus implies $\widehat{\delta} \widehat{d} v = 0$ and $\widehat{d} \widehat{\delta} v = 0$, hence $\widehat{\delta} \widehat{d} \widetilde{v}$ and $\widehat{d} \widehat{\delta} \widetilde{v}$ are supported in $\widetilde{X} \setminus X$. (This argument shows the uniqueness of the decomposition (3.42).) Therefore \widetilde{u} is a solution of $\widetilde{\square} \widetilde{u} = 0$ which is supported in $\widetilde{X} \setminus X$. By unique continuation at infinity on the asymptotically de Sitter side $\widetilde{X} \setminus X$ of \widetilde{X} , this implies $\widetilde{u} \equiv 0$, as claimed.

Now if $u \in \mathcal{C}_{(0)}^\infty$ is given, extend it arbitrarily to $\widetilde{u} \in \mathcal{C}^\infty(\widetilde{X})$, apply (3.43) and restrict both summands back to X . This establishes (3.42). \square

Remark 3.13. The decomposition (3.42) does *not* hold if we replace $\widehat{\square}$ in (3.42) by $\widehat{d} + \widehat{\delta}$. Indeed, if it did hold, this would say that $\widehat{\square} u = 0$ implies $(\widehat{d} + \widehat{\delta})u = 0$, since $(\widehat{d} + \widehat{\delta})u$ lies both in $\ker(\widehat{d} + \widehat{\delta})$ and $\text{ran}(\widehat{d} + \widehat{\delta})$ in this case. Since certainly $(\widehat{d} + \widehat{\delta})u = 0$ conversely implies $\widehat{\square} u = 0$, this would mean that $\ker \widehat{\square} = \ker(\widehat{d} + \widehat{\delta})$. Now by Lemmas 2.1 and 2.2, this in turn would give $\ker \widetilde{\square} = \ker(\widetilde{d} + \widetilde{\delta})$. Now since $\widetilde{\square}$ and $\widetilde{d} + \widetilde{\delta}$ are Fredholm with index 0, we could further deduce $\ker \widetilde{\square}^* = \ker(\widetilde{d} + \widetilde{\delta})^*$, where the adjoints act on the space $\dot{\mathcal{C}}^{-\infty}(\widetilde{X})$ of supported distributions at the (artificial) Cauchy hypersurface $\partial \widetilde{X}$, see [31, Appendix B]. Since we have $\ker(\widetilde{d} + \widetilde{\delta})^* \subset \ker \widetilde{\square}^*$ unconditionally, we can show the absurdity of this last equality by exhibiting an element u in $\ker \widetilde{\square}^*$ which does not lie in $\ker(\widetilde{d} + \widetilde{\delta})^*$. This however is easy: Just let $u = 1_X$ be the characteristic function of X . Then from (3.14) and (3.15), we

see that $(\tilde{d} + \tilde{\delta})u = \tilde{d}u$ is a non-zero delta distribution supported at Y which is annihilated by $\tilde{\delta}$.

This argument shows that we always have $\ker \hat{\square} \supsetneq \ker(\hat{d} + \hat{\delta})$. It is possible though that $\mathcal{H}^k = \mathcal{K}^k$ for *some* form degrees k (but this must fail for some value of k). For instance, this holds for $k = 0$ by Corollary 3.9. We will give a more general statement below, see in particular Remark 3.18.

We now define a complex whose cohomology we will relate to the spaces \mathcal{H}^k and \mathcal{K}^k : The space $\mathcal{C}_{(0)}^\infty \cap L^2(\alpha|dh|)$ of smooth forms $u = \mathcal{C}\tilde{u}$ with $\tilde{u}_{NT}|_Y = 0$ has a grading corresponding to form degrees, thus

$$\mathcal{D} := \mathcal{C}_{(0)}^\infty \cap L^2(\alpha|dh|) = \bigoplus_{k=0}^n \mathcal{D}^k.$$

Since in the above notation $u \in L^2(\alpha|dh|)$ (and thus $\tilde{u}_{NT}|_Y = 0$) is equivalent to $\tilde{u}_{NT} \in \alpha^2 \mathcal{C}^\infty(\overline{X}_{\text{even}}; \Lambda Y)$ near Y , one can easily check using (3.14) that \hat{d} acts on $\mathcal{C}_{(0)}^\infty \cap L^2(\alpha|dh|)$. We can then define the complex

$$0 \rightarrow \mathcal{D}^0 \xrightarrow{\hat{d}} \mathcal{D}^1 \rightarrow \dots \xrightarrow{\hat{d}} \mathcal{D}^n \rightarrow 0.$$

We denote its cohomology by

$$\mathcal{H}_{L^2, \text{dR}}^k = \ker(\hat{d}: \mathcal{D}^k \rightarrow \mathcal{D}^{k+1}) / \text{ran}(\hat{d}: \mathcal{D}^{k-1} \rightarrow \mathcal{D}^k). \quad (3.44)$$

There is a natural map from $\mathcal{H}_{L^2, \text{dR}}^k$ into \mathcal{H}^k :

Lemma 3.14. *Every cohomology class $[u] \in \mathcal{H}_{L^2, \text{dR}}^k$ has a unique representative $u' \in \mathcal{H}^k$, and the map $i: [u] \mapsto u'$ is injective.*

Proof. Let $[u] \in \mathcal{H}_{L^2, \text{dR}}^k$, hence $\hat{d}u = 0$ and, writing $u = \mathcal{C}\tilde{u}$, $\tilde{u}_{NT}|_Y = 0$. We first show the existence of a representative, i.e. an element $u - \hat{d}v$ with $v \in \mathcal{D}$, which is annihilated by $\hat{\delta}$. (Since it is clearly annihilated by \hat{d} , this means $u - \hat{d}v \in \mathcal{H}^k$.) That is, we need to solve the equation $\hat{\delta}\hat{d}v = \hat{\delta}u$ with $v \in \mathcal{D}$. To achieve this, we use Lemma 3.12 to write

$$u = u_1 + (\hat{d}\hat{\delta} + \hat{\delta}\hat{d})u_2, \quad u_1 \in \ker \hat{\square}.$$

By our assumption on u and Corollary 3.10, u and u_1 are annihilated by $\hat{\delta}\hat{d}$, giving $\hat{\delta}\hat{d}\hat{\delta}\hat{d}u_2 = 0$. By Lemma 3.11, this implies $\hat{\delta}\hat{d}u_2 = 0$, hence

$$u = u_1 + \hat{d}\hat{\delta}u_2. \quad (3.45)$$

Applying $\hat{d}\hat{\delta}$, we obtain

$$\hat{d}\hat{\delta}\hat{d}\hat{\delta}u_2 = \hat{d}\hat{\delta}u \in L^2. \quad (3.46)$$

Now writing $u_2 = \mathcal{C}\tilde{u}_2$, and noting that for any $w = \mathcal{C}\tilde{w} \in \mathcal{C}_{(0)}^\infty$, $(\mathcal{C}^{-1}\hat{d}\mathcal{C}\tilde{w})_{NT}|_Y = -d_Y\tilde{w}_{NT}|_Y$ as well as $(\mathcal{C}^{-1}\hat{\delta}\mathcal{C}\tilde{w})_{NT}|_Y = \delta_Y\tilde{w}_{NT}|_Y$ by (3.14) and (3.15), the (NT) component of \mathcal{C}^{-1} times equation (3.46) reads $d_Y\delta_Y d_Y\delta_Y\tilde{u}_{2,NT}|_Y = 0$, which yields $\delta_Y\tilde{u}_{2,NT}|_Y = 0$. As a consequence of this, $v := \hat{d}u_2 \in L^2$ and therefore $\hat{d}\hat{\delta}u_2 \in L^2$. Hence (3.45) gives $u_1 \in L^2$; by Corollary 3.9 then, $u_1 \in \mathcal{H}$, in particular u_1 is annihilated by $\hat{\delta}$. Therefore, applying $\hat{\delta}$ to (3.45) yields $\hat{\delta}(u - \hat{d}v) = 0$, as desired.

Next, we show that the representative is unique: Thus, suppose $u - \hat{d}v_1, u - \hat{d}v_2 \in \mathcal{H}^k$ with $u, v_1, v_2 \in \mathcal{D}$, then with $v = v_1 - v_2 \in \mathcal{D}$, we have $\hat{d}v \in \mathcal{H}^k$, thus $\hat{\delta}\hat{d}v = 0$,

and by Corollary 3.9, we obtain $\widehat{d}v = 0$. Therefore, $u - \widehat{d}v_1 = u - \widehat{d}v_2$, establishing uniqueness, which in particular shows that the map i is well-defined.

Finally, we show the injectivity of i : Suppose $u \in \mathcal{D}$ satisfies $\widehat{d}u = 0$. There exists an element $v \in \mathcal{D}$ such that $u - \widehat{d}v \in \mathcal{H}^k$. Now if $i[u] = 0$, this precisely means that $u - \widehat{d}v = 0$; but then $[u] = [\widehat{d}v] = 0$ in $\mathcal{H}_{L^2, \text{dR}}^k$. \square

From the definition of the space \mathcal{D} , it is clear that $u \in \mathcal{H}^k$ lies in the image of i if and only if $u \in L^2$, i.e. if and only if $r(u) = 0$, where r is the map

$$r: \mathcal{C}_{(0)}^\infty \rightarrow \mathcal{C}^\infty(Y; \Lambda Y), \quad u = \mathcal{C}\tilde{u} \mapsto \tilde{u}_{NT}|_Y. \quad (3.47)$$

Thus, r extracts the singular part of u and thereby measures the failure of a given form $u \in \mathcal{C}_{(0)}^\infty$ to lie in \mathcal{D} . Observe that if $u = \mathcal{C}\tilde{u} \in \mathcal{H}^k$, then $d_Y \tilde{u}_{NT}|_Y = 0$ and $\delta_Y \tilde{u}_{NT}|_Y = 0$, i.e. $r(u)$ is a harmonic form on Y . Since the space $\ker(\Delta_{Y, k-1})$ of harmonic forms on the closed manifold Y is isomorphic to the cohomology group $H^{k-1}(Y)$ by standard Hodge theory, we thus obtain:

Proposition 3.15. *The sequence*

$$0 \rightarrow \mathcal{H}_{L^2, \text{dR}}^k \xrightarrow{i} \mathcal{H}^k \xrightarrow{r} H^{k-1}(Y) \quad (3.48)$$

is exact. Here, i is the map defined in Lemma 3.14, and r is the restriction map (3.47) (composed with the identification $\ker(\Delta_{Y, k-1}) \cong H^{k-1}(Y)$). Moreover, the map $i: \mathcal{H}_{L^2, \text{dR}}^k \rightarrow \mathcal{H}^k \cap \mathcal{D}$ is an isomorphism with inverse $\mathcal{H}^k \cap \mathcal{D} \ni u \mapsto [u] \in \mathcal{H}_{L^2, \text{dR}}^k$.

Proof. We only need to check the last claim. If $u \in \mathcal{H}^k \cap \mathcal{D}$, then $[u]$ does define a cohomology class in $\mathcal{H}_{L^2, \text{dR}}^k$, and $i([u])$ is the unique representative of $[u]$ which lies in \mathcal{H}^k . Since u itself is such a representative, we must have $i([u]) = u$. For the converse, we note that for any $[u] \in \mathcal{H}_{L^2, \text{dR}}^k$ we have $i([u]) = u - \widehat{d}v$ for some $v \in \mathcal{D}$, hence $[i([u])] = [u - \widehat{d}v] = [u]$. \square

We can make a stronger statement: If we merely have $u \in \ker \widehat{\square}$, then the proof of Proposition 3.8 shows that $\widehat{d}u, \widehat{\delta}u \in L^2$, hence $r(u)$ is harmonic.

Proposition 3.16. *We have a short exact sequence*

$$0 \rightarrow \mathcal{H}_{L^2, \text{dR}}^k \xrightarrow{i} \mathcal{K}^k \xrightarrow{r} H^{k-1}(Y) \rightarrow 0, \quad (3.49)$$

where the first map is i defined in Lemma 3.14 (composed with the inclusion $\mathcal{H}^k \hookrightarrow \mathcal{K}^k$), and the second map is the restriction r , defined in (3.47) (composed with the identification $\ker(\Delta_{Y, k-1}) \cong H^{k-1}(Y)$).

Proof. The second map is well-defined by the comment preceding the statement of the proposition. Since the range of $\mathcal{H}_{L^2, \text{dR}}^k$ in \mathcal{K}^k consists of L^2 forms, we have $r \circ i = 0$. Moreover, if $u \in \ker r$, then u is an L^2 element of $\ker \widehat{\square}$, thus $u \in \mathcal{H}^k$ by Corollary 3.9. By the remark following the proof of Lemma 3.14, therefore $u \in \text{ran } i$.

It remains to show the surjectivity of r : Thus, let $w \in \ker(\Delta_{Y, k-1})$, and let $u' = \mathcal{C}\tilde{u}' \in \mathcal{C}_{(0)}^\infty$ be any extension of w , i.e. $\tilde{u}'_{NT}|_Y = w$. Then $(\widehat{d} + \widehat{\delta})u' \in \mathcal{D}$, since its (NT) component vanishes, and thus $\widehat{\square}u' \in \mathcal{D}$. Writing $u' = u_1 + \widehat{\square}u_2$ with $u_1 \in \ker \widehat{\square}$, we conclude that $\widehat{\square}u' = \widehat{\square}^2u_2$; taking the (NT) component of this equation gives $0 = \Delta_Y^2 \tilde{u}_{2, NT}|_Y$ (where we write $u_2 = \mathcal{C}\tilde{u}_2$ as usual), hence

$d_Y \tilde{u}_{2,NT}|_Y = 0$ and $\delta_Y \tilde{u}_{2,NT}|_Y = 0$. But then $\widehat{\square} u_2 \in L^2$. Therefore, $w = r(u') = r(u_1 + \widehat{\square} u_2) = r(u_1)$. Since the degree k part of u_1 lies in \mathcal{K}^k by the definition of u_1 , we are done. \square

Remark 3.17. Remark 3.13, which states that $\mathcal{H}^k \subsetneq \mathcal{K}^k$ for some values of k , implies in particular that the last map of (3.48) is not always onto.

Remark 3.18. Since $\dim Y = n - 2$, we have $H^{k-1}(Y) = 0$ for $k = 0$ and $k = n$. Hence, for these extreme values of k , Propositions 3.15 and 3.16 show $\mathcal{H}^k = \mathcal{K}^k \cong \mathcal{H}_{L^2, \text{dR}}^k$, and this holds more generally for all k for which $H^{k-1}(Y) = 0$.

The spaces $\mathcal{H}_{L^2, \text{dR}}^k$ are related to standard cohomology groups associated with the manifold with boundary \overline{X} : First, notice that elements of the space $\mathcal{D} = \mathcal{C}_{(0)}^\infty \cap L^2$ are not subject to any matching condition on singular terms, simply because the singular term ($\tilde{u}_{NT}|_Y$ in the notation used above) vanishes. This means that we can split \mathcal{D} into tangential and normal forms, $\mathcal{D} = \mathcal{D}_T \oplus \mathcal{D}_N$,¹¹ where \mathcal{D}_T consists of all $u_T \in \mathcal{C}^\infty(\overline{X}; \Lambda \overline{X})$ which are of the form

$$u_T = \begin{pmatrix} u_{TT} \\ \alpha u_{TN} \end{pmatrix}, \quad u_{TT}, u_{TN} \in \mathcal{C}^\infty(\overline{X}_{\text{even}}; \Lambda Y),$$

near Y . Thus, elements $u_T \in \mathcal{D}_T$ are forms of the type $u_T = u_{TT} + d\alpha \wedge \alpha u_{TN} = u_{TT} + \frac{1}{2}d\mu \wedge u_{TN}$ with u_{TT}, u_{TN} smooth ΛY -valued forms on $\overline{X}_{\text{even}}$; hence, we simply have $\mathcal{D}_T = \mathcal{C}^\infty(\overline{X}_{\text{even}}; \Lambda \overline{X}_{\text{even}})$. Likewise, \mathcal{D}_N consists of all $u_N \in \mathcal{C}^\infty(\overline{X}; \Lambda \overline{X})$ which are of the form

$$u_N = \begin{pmatrix} \alpha u_{NT} \\ u_{NN} \end{pmatrix}, \quad u_{NT}, u_{NN} \in \mathcal{C}^\infty(\overline{X}_{\text{even}}; \Lambda Y),$$

near Y . Thus, elements $u_N \in \mathcal{D}_N$ are forms of the type $\alpha u_N = \mu u_{NT} + \frac{1}{2}d\mu \wedge u_{NN}$; therefore, $\alpha \mathcal{D}_N = \mathcal{C}_R^\infty(\overline{X}_{\text{even}}; \Lambda \overline{X}_{\text{even}}) := \{u \in \mathcal{C}^\infty(\overline{X}_{\text{even}}; \Lambda \overline{X}_{\text{even}}) : j^* u = 0\}$, where $j : \partial \overline{X}_{\text{even}} \hookrightarrow \overline{X}_{\text{even}}$ is the inclusion.

Since the differential \widehat{d} on \mathcal{D} acts as $d_X \oplus (-\alpha^{-1} d_X \alpha)$ on $\mathcal{D}_T \oplus \mathcal{D}_N$, the cohomology of the complex $(\mathcal{D}, \widehat{d})$ in degree k is the direct sum of the cohomology of (\mathcal{D}_T, d_X) in degree k and of $(\alpha \mathcal{D}_N, d_X)$ in degree $(k-1)$. Since we identified \mathcal{D}_T as simply the space of smooth forms on $\overline{X}_{\text{even}}$, the cohomology of (\mathcal{D}_T, d_X) in degree k equals the absolute cohomology $H^k(\overline{X}_{\text{even}}) \cong H^k(\overline{X})$.¹² Moreover, since \mathcal{D}_N is the space of smooth forms on $\overline{X}_{\text{even}}$ which vanish at the boundary in the precise sense described above, the cohomology of $(\alpha \mathcal{D}_N, d_X)$ in degree k equals the relative cohomology $H^k(\overline{X}_{\text{even}}; \partial \overline{X}_{\text{even}}) \cong H^k(\overline{X}; \partial \overline{X})$ (see e.g. [49, §5.9]). In summary:

Proposition 3.19. *With $\mathcal{H}_{L^2, \text{dR}}^k$ defined in (3.44), there is a canonical isomorphism*

$$\mathcal{H}_{L^2, \text{dR}}^k \cong H^k(\overline{X}) \oplus H^{k-1}(\overline{X}, \partial \overline{X}). \quad (3.50)$$

Let us summarize the results obtained in the previous sections:

Theorem 3.20. *The only resonance of $d + \delta$ in $\text{Im } \sigma \geq 0$ is $\sigma = 0$, and 0 is a simple resonance. Zero resonant states of the extended operator $(d + \delta)$ on \widehat{M} are uniquely determined by their restriction to X , and the space \mathcal{H} of these resonant states on*

¹¹Thus, elements $(u_T, u_N) \in \mathcal{D}_T \oplus \mathcal{D}_N$ are identified with $u_T + \alpha dt \wedge u_N \in \mathcal{D}$.

¹²We use that $\overline{X}_{\text{even}}$ is diffeomorphic to \overline{X} , with diffeomorphism given by gluing the map $\alpha^2 \mapsto \alpha$ near Y to the identity map away from Y .

X is equal to $\ker_{\mathcal{C}_{(0)}^\infty} \widehat{d}(0) \cap \ker_{\mathcal{C}_{(0)}^\infty} \widehat{\delta}(0)$. Also, resonant states on \widetilde{X} are elements of $\ker \widetilde{d}(0) \cap \ker \widetilde{\delta}(0)$. Using the grading $\mathcal{H} = \bigoplus_{k=0}^n \mathcal{H}^k$ of \mathcal{H} by form degrees, there is a canonical exact sequence

$$0 \rightarrow H^k(\overline{X}) \oplus H^{k-1}(\overline{X}, \partial\overline{X}) \rightarrow \mathcal{H}^k \rightarrow H^{k-1}(\partial\overline{X}), \quad (3.51)$$

where the first map is the composition of the isomorphism (3.50) with the map i defined in Lemma 3.14, and the second map is the composition of the map r defined in (3.47) with the isomorphism $\ker(\Delta_{\partial\overline{X}, k-1}) \cong H^{k-1}(\partial\overline{X})$.

Furthermore, the only resonance of \square_g in $\text{Im } \sigma \geq 0$ is $\sigma = 0$. Zero resonant states¹³ of the extended operator $(\square_g \text{ on } \widetilde{M})$ are uniquely determined by their restriction to X . The space $\mathcal{K} = \bigoplus_{k=0}^n \mathcal{K}^k \subset \mathcal{C}_{(0)}^\infty$ of these resonant states on X , graded by form degree, satisfying $\mathcal{K}^k \supset \mathcal{H}^k$, fits into the short exact sequence

$$0 \rightarrow H^k(\overline{X}) \oplus H^{k-1}(\overline{X}, \partial\overline{X}) \rightarrow \mathcal{K}^k \rightarrow H^{k-1}(\partial\overline{X}) \rightarrow 0, \quad (3.52)$$

with maps as above. We moreover have

$$\mathcal{K}^k \cap L^2 = \mathcal{H}^k \cap L^2 \cong H^k(\overline{X}) \oplus H^{k-1}(\overline{X}, \partial\overline{X})$$

where $L^2 = L^2(X, \alpha|dh|; H \oplus H)$. More precisely then, the summand $H^k(\overline{X})$ in (3.51) and (3.52) corresponds to the tangential components (in the decomposition (3.1)) of elements of $\mathcal{H}^k \cap L^2$, and the summand $H^{k-1}(\overline{X}, \partial\overline{X})$ to the normal components.

Lastly, the Hodge star operator on M induces isomorphisms $\star: \mathcal{H}^k \xrightarrow{\cong} \mathcal{H}^{n-k}$ and $\star: \mathcal{K}^k \xrightarrow{\cong} \mathcal{K}^{n-k}$, $k = 0, \dots, n$.

Proof. We prove the statement about resonant states for $d + \delta$ on the extended space \widetilde{M} : Thus, if $\widetilde{u} \in \ker(\widetilde{d}(0) + \widetilde{\delta}(0))$, then the restriction of \widetilde{u} to X lies in $\ker \widehat{d}(0) \cap \ker \widehat{\delta}(0)$, therefore $\widehat{d}(0)\widetilde{u} = -\widehat{\delta}(0)\widetilde{u}$ is supported in $\widetilde{X} \setminus X$; but then $\widetilde{\square}(0)(\widehat{d}(0)\widetilde{u}) = \widehat{d}(0)\widehat{\delta}(0)\widehat{d}(0)\widetilde{u} = 0$ and the asymptotically de Sitter nature of $\widetilde{X} \setminus X$ implies $\widehat{d}(0)\widetilde{u} \equiv 0$, hence also $\widehat{\delta}(0)\widetilde{u} \equiv 0$, as claimed.

The only remaining part of the statement that has not yet been proved is the last: Viewing $u \in \mathcal{H}^k$ as a t -independent k -form on $M = \mathbb{R}_t \times X$ (with the metric (1.1)), we have $(d + \delta)u = 0$, and for any t -independent k -form u on M , we have that $(d + \delta)u = 0$ implies $u \in \mathcal{H}^k$, where we view the t -independent form as a form on X valued in the form bundle of M , as explained in Section 2. Then $u \in \mathcal{H}^k$ is equivalent to $du = 0$, $\delta u = 0$, which in turn is equivalent to $\delta(\star u) = 0$, $d(\star u) = 0$, and thus $\star u \in \mathcal{H}^{n-k}$. The proof for the spaces \mathcal{K}^k is the same and uses $\star\square = \square\star$. \square

This in particular proves Theorem 2.

4. RESULTS FOR STATIC DE SITTER AND SCHWARZSCHILD-DE SITTER SPACETIMES

We now supplement the results obtained in the previous section by high energy estimates for the inverse normal operator family and deduce expansions and decay for solutions to Maxwell's equations as well as for more general linear waves on de Sitter and Schwarzschild-de Sitter backgrounds. The rather detailed description of

¹³More precisely, we mean elements of $\ker \widetilde{\square}(0)$; the latter space equals the space of zero resonant states if the zero resonance is simple.

asymptotics in the Schwarzschild-de Sitter setting will be essential in our discussion of Kerr-de Sitter space in Section 5.

4.1. de Sitter space. De Sitter space is the hyperboloid $\{|x|^2 - \tilde{t}^2 = 1\}$ in $(n+1)$ -dimensional Minkowski space, equipped with the induced Lorentzian metric. Introducing $\tau = \tilde{t}^{-1}$ in $\tilde{t} \geq 1$ and adding the boundary at future infinity $\tau = 0$ to the spacetime, we obtain the bordified space $N = [0, 1)_\tau \times Z$ with $Z = \mathbb{S}^{n-1}$, and the metric has the form

$$g^0 = \tau^{-2} \bar{g}, \quad \bar{g} = d\tau^2 - h^0(\tau, x, dx),$$

with h^0 even in τ , i.e. h^0 is a metric on Z which depends smoothly on τ^2 ; see Vasy [52, §4] for details. Thus, g^0 is a 0-metric in the sense of Mazzeo and Melrose [39]. Fixing a point p at future infinity, the static model of de Sitter space, denoted M , is the interior of the backward light cone from p .¹⁴ We introduce *static coordinates* on M , denoted $(t, x) \in \mathbb{R} \times X$, where $X = B_1 \subset \mathbb{R}^{n-1}$ is the open unit ball in \mathbb{R}^{n-1} and $x \in \mathbb{R}^{n-1}$ are the standard coordinates on \mathbb{R}^{n-1} , with respect to which the induced metric on M is given by

$$g = \alpha^2 dt^2 - h, \quad \alpha = (1 - |x|^2)^{1/2},$$

$$h = dx^2 + \frac{1}{1 - |x|^2} (x \cdot dx)^2 = \alpha^{-2} dr^2 + r^2 d\omega^2,$$

using polar coordinates (r, ω) on \mathbb{R}_x^{n-1} near $r = 1$, and denoting the round metric on the unit sphere \mathbb{S}^{n-2} by $d\omega^2$. We compactify X to the closed unit ball $\bar{X}_{\text{even}} = \bar{B}_1 \subset \mathbb{R}^{n-1}$, and denote by \bar{X} the space which is \bar{X}_{even} topologically, but with α added to the smooth structure. In order to see that the metric g fits into the framework of Theorem 3.20, note that $dr = -\alpha r^{-1} d\alpha$, so

$$h = r^{-2} d\alpha^2 + r^2 d\omega^2,$$

and $r = (1 - \alpha^2)^{1/2}$, thus h is an even metric on the space \bar{X} and has the form (1.2) with $\beta = 1$. Using Theorem 3.20, we can now easily compute the spaces of resonances:

Theorem 4.1. *On an n -dimensional static de Sitter spacetime, $n \geq 4$, the spaces of resonances of \square and $d + \delta$ are*

$$\mathcal{K}^0 = \mathcal{H}^0 = \langle 1 \rangle, \quad \mathcal{K}^n = \mathcal{H}^n = \langle r^{n-2} dt \wedge dr \wedge \omega \rangle,$$

where ω denotes the volume form on the round sphere \mathbb{S}^{n-2} . Furthermore,

$$\mathcal{K}^1 = \langle -\alpha^{-2} r dr + \alpha^{-1} dt \rangle, \mathcal{H}^1 = 0, \quad \mathcal{K}^{n-1} = \langle \star(-\alpha^{-2} r dr + \alpha^{-1} dt) \rangle, \mathcal{H}^{n-1} = 0,$$

$$\mathcal{K}^k = \mathcal{H}^k = 0, \quad k = 2, \dots, n-2.$$

In particular, on 4-dimensional static de Sitter space, if u is a solution of $(d + \delta)u = 0$ with smooth initial data, then the degree 0 component of u decays exponentially to a constant, the degree 1, 2 and 3 components decay exponentially to 0, and the degree 4 component decays exponentially to a constant multiple of the volume form. Analogous statements hold on any n -dimensional static de Sitter space, $n \geq 5$.

¹⁴Since g^0 and the metric \bar{g} , which is smooth down to $\tau = 0$, are conformally related, the images of null-geodesics for both metrics agree.

Proof. We compute the cohomological data that appear in (3.51) and (3.52) using $\overline{X} \cong \overline{B_1}$ and $\partial\overline{X} \cong \mathbb{S}^{n-2}$:

$$\begin{aligned} \dim H^{k-1}(\partial\overline{X}) &= \begin{cases} 0, & k = 0, 2, \dots, n-2, n, \\ 1, & k = 1, n-1 \end{cases} \\ \dim H^k(\overline{X}) &= \begin{cases} 1, & k = 0 \\ 0, & 1 \leq k \leq n, \end{cases} \\ \dim H^{k-1}(\overline{X}, \partial\overline{X}) &= \begin{cases} 0, & 0 \leq k \leq n-1 \\ 1, & k = n. \end{cases} \end{aligned}$$

Thus, we immediately deduce

$$\begin{aligned} \dim \mathcal{K}^0 &= \dim \mathcal{K}^1 = \dim \mathcal{K}^{n-1} = \dim \mathcal{K}^n = 1, \quad \dim \mathcal{K}^k = 0, \quad 2 \leq k \leq n-2, \\ \dim \mathcal{H}^0 &= \dim \mathcal{H}^n = 1, \quad \dim \mathcal{H}^k = 0, \quad 2 \leq k \leq n-2. \end{aligned}$$

Now, since $d + \delta$ annihilates constants, we find $1 \in \mathcal{K}^0 = \mathcal{H}^0$ and $\star 1 \in \mathcal{K}^n = \mathcal{H}^n$, which in view of the 1-dimensionality of these spaces already concludes their computation.

In order to compute \mathcal{K}^1 , notice that we have $\mathcal{K}^1 \cong H^0(\partial\overline{X})$ from (3.52), thus an element u spanning \mathcal{K}^1 has non-trivial singular components at $\alpha = 0$. One is led to the guess $u = \alpha^{-1} d\alpha + \alpha^{-1} dt = -\alpha^{-2} r dr + \alpha^{-1} dt$, which is indeed annihilated by \square ; we will give full details for this computation in the next section when discussing Schwarzschild-de Sitter spacetimes, which in the case of vanishing black hole mass are static de Sitter spacetimes, with a point removed, see in particular the calculations following (4.10); but since u as defined above is smooth at $r = 0$, we obtain $\square u = 0$ at $r = 0$ as well by continuity. Since \mathcal{K}^1 is 1-dimensional, we therefore deduce $\mathcal{K}^1 = \langle u \rangle$. One can then check that $(d + \delta)u \neq 0$, and this implies $\mathcal{H}^1 = 0$. The corresponding statements for \mathcal{K}^{n-1} and \mathcal{H}^{n-1} are immediate consequences of this and the fact that the Hodge star operator induces isomorphisms $\mathcal{H}^1 \cong \mathcal{H}^{n-1}$ and $\mathcal{K}^1 \cong \mathcal{K}^{n-1}$.

The high energy estimates for $d + \delta$ required to deduce asymptotic expansions for solutions of $(d + \delta)u = 0$ follow from those of its square \square , which is principally scalar and fits directly into the framework described in [52, §2-4]. \square

By studying the space of dual resonant states, one can in fact easily show that the 0-resonance of \square is simple and thus deduce exponential decay of smooth solutions to $\square u = 0$ to an element of \mathcal{K}^k in all form degrees $k = 0, \dots, n$. We give details in the next section on Schwarzschild-de Sitter space.

In the present de Sitter setting, one can deduce asymptotics very easily in a different manner using the global de Sitter space picture, by analyzing indicial operators in the 0-calculus: Concretely, we write differential k -forms (by which we mean smooth sections of the k -th exterior power of the 0-cotangent bundle of N) as

$$u = \tau^{-k} u_T + \frac{d\tau}{\tau} \wedge \tau^{1-k} u_N, \quad (4.1)$$

where u_T and u_N are smooth forms on Z of form degrees k and $(k-1)$, respectively. One readily computes the differential d_k acting on k -forms to be

$$d_k = \begin{pmatrix} \tau d_Z & 0 \\ -k + \tau \partial_\tau & -\tau d_Z \end{pmatrix}.$$

Furthermore, by the choice of basis in (4.1), the inner product on k -forms induced by g^0 is given by

$$G_k^0 = \begin{pmatrix} (-1)^k H_k^0 & 0 \\ 0 & (-1)^{k-1} H_{k-1}^0 \end{pmatrix}.$$

Using that the volume density is $|dg^0| = \tau^{-n} d\tau |dh^0|$, we compute the codifferential δ_k acting on k -forms to be

$$\delta_k = \begin{pmatrix} -\tau\delta_Z & -(k-1) + \tau^{n-1}\tau\partial_\tau^*\tau^{1-n} \\ 0 & \tau\delta_Z \end{pmatrix} = \begin{pmatrix} -\tau\delta_Z & n-k - \tau\partial_\tau + \mathcal{O}_{C^\infty(N)}(\tau) \\ 0 & \tau\delta_Z \end{pmatrix},$$

where ∂_τ^* is the $L^2(N, |d\bar{g}|)$ -adjoint (suppressing the bundles in the notation) of ∂_τ , and we use the even-ness of g^0 in the second step to deduce $\partial_\tau^* = -\partial_\tau + \mathcal{O}_{C^\infty(N)}(\tau)$. Therefore, the indicial roots of $d + \delta$ on the degree k -part of the form bundle are k and $n-k$.

Next, for $0 \leq k \leq n$, we compute the Hodge d'Alembertian:¹⁵

$$\begin{aligned} \square_k &= d_{k-1}\delta_k + \delta_{k+1}d_k \\ &= \begin{pmatrix} -\tau d_Z \tau \delta_Z - \tau \delta_Z \tau d_Z - P_k & \tau d_Z \\ -\tau \delta_Z & -\tau d_Z \tau \delta_Z - \tau \delta_Z \tau d_Z - P_{k-1} \end{pmatrix} + \mathcal{O}_{\text{Diff}_0^1}(\tau) \end{aligned}$$

where $P_k = (\tau\partial_\tau)^2 - (n-1)\tau\partial_\tau + k(n-k-1)$. Thus, the indicial polynomial of \square_k is

$$I(\square_k)(s) = \begin{pmatrix} s^2 - (n-1)s + k(n-k-1) & 0 \\ 0 & s^2 - (n-1)s + (k-1)(n-k) \end{pmatrix}.$$

On tangential forms, the indicial roots of \square_k are therefore $k, n-1-k$, and on normal forms, they are $k-1, n-k$. We thus have:

form degree	0	1	$2 \leq k \leq n-2$	$n-1$	n
tgt. ind. roots	$0, n-1$	$1, n-2$	$k, n-1-k$	$0, n-1$	—
norm. ind. roots	—	$0, n-1$	$k-1, n-k$	$1, n-2$	$0, n-1$

Hence in particular, all roots are ≥ 0 , and 0 is never a double root. Thus, the arguments of [50] (which are in the scalar setting, but work in the current setting as well with only minor modifications) show that solutions u to the wave equation on differential k -forms on N with smooth initial data at $\tau = \tau_0 > 0$ decay exponentially (in $-\log \tau$) if 0 is not an indicial root, and decay to a stationary state if 0 is an indicial root.¹⁶ Explicitly, scalar waves decay to a smooth function on Z , 1-form waves decay to an element of $\frac{d\tau}{\tau} \mathcal{C}^\infty(Z)$, k -form waves decay exponentially to 0 for $2 \leq k \leq n-2$, $(n-1)$ -form waves decay to an element of $\mathcal{C}^\infty(Z; \Lambda^{n-1}Z)$, and n -form waves finally decay to an element of $\frac{d\tau}{\tau} \wedge \mathcal{C}^\infty(Z; \Lambda^{n-1}Z)$.

Since the static model of de Sitter space arises by blowing up a point p at future infinity of compactified de Sitter space and considering the backward light cone from p , we can find the resonant states for the static model by simply finding the space of restrictions to p of the asymptotic states described above; but since the fibers of $\Lambda^0(Z)$ and $\Lambda^{n-1}(Z)$ are 1-dimensional, hence we have reproved Theorem 4.1.

¹⁵We deal with the cases $k=0$ and $k=n$ simultaneously with $1 \leq k \leq n-1$ by implicitly assuming that for $k=0$, only the $(1,1)$ -part of this operator acts on 0-forms, and for $k=n$, only the $(2,2)$ -part acts on n -forms.

¹⁶Of course, since we know all indicial roots, we could be much more precise in describing the asymptotics, but we only focus on the 0-resonance here.

We point out that if one wants to analyze differential form-valued waves or solutions to Maxwell's equations on Schwarzschild-de Sitter space, there is no global picture (in the sense of a 0-differential problem) as in the de Sitter case. Thus, the direct approach outlined in the proof of Theorem 4.1 is the only possible one in this case, and it is very instructive as it shows even more clearly how the cohomological interpretation of the space of zero resonant states can be used very effectively.

4.2. Schwarzschild-de Sitter space. The computation of resonant states for Schwarzschild-de Sitter spacetimes of any dimension is no more difficult than the computation in 4 dimensions, thus we directly treat the general case of $n \geq 4$ space-time dimensions. Recall that the metric of n -dimensional Schwarzschild-de Sitter space $M = \mathbb{R}_t \times X$, $X = (r_-, r_+)_r \times \mathbb{S}_\omega^{n-2}$, with r_\pm defined below, is given by

$$g = \mu dt^2 - (\mu^{-1} dr^2 + r^2 d\omega^2),$$

where $d\omega^2$ is the round metric on the sphere \mathbb{S}^{n-2} , and $\mu = 1 - \frac{2M_\bullet}{r^{n-3}} - \lambda r^2$, $\lambda = \frac{2\Lambda}{(n-2)(n-1)}$, where the black hole mass M_\bullet and the cosmological constant Λ are positive. We assume that

$$M_\bullet^2 \lambda^{n-3} < \frac{(n-3)^{n-3}}{(n-1)^{n-1}}, \quad (4.2)$$

which guarantees that μ has two unique positive roots $0 < r_- < r_+$. Indeed, let $\tilde{\mu} = r^{-2}\mu = r^{-2} - 2M_\bullet r^{1-n} - \lambda$. Then $\tilde{\mu}' = -2r^{-n}(r^{n-3} - (n-1)M_\bullet)$ has a unique positive root $r_p = [(n-1)M_\bullet]^{1/(n-3)}$, $\tilde{\mu}'(r) > 0$ for $r \in (0, r_p)$ and $\tilde{\mu}'(r) < 0$ for $r > r_p$; moreover, $\tilde{\mu}(r) < 0$ for $r > 0$ small and $\tilde{\mu}(r) \rightarrow -\lambda < 0$ as $r \rightarrow \infty$, thus the existence of the roots $0 < r_- < r_+$ of $\tilde{\mu}$ is equivalent to the requirement $\tilde{\mu}(r_p) = \frac{n-3}{n-1}r_p^{-2} - \lambda > 0$, which leads precisely to the inequality (4.2).

Define $\alpha = \mu^{1/2}$, thus $d\alpha = \frac{1}{2}\mu'\alpha^{-1}dr$, and

$$\beta_\pm := \mp \frac{2}{\mu'(r_\pm)} > 0,$$

then the metric g can be written as

$$g = \alpha^2 dt^2 - h, \quad h = \tilde{\beta}_\pm^2 d\alpha^2 + r^2 d\omega^2,$$

where $\tilde{\beta}_\pm = \mp 2/\mu'(r)$. Thus, if we let $\bar{X}_{\text{even}} = [r_-, r_+]_r \times \mathbb{S}_\omega^{n-2}$ with the standard smooth structure, then $\tilde{\beta}_\pm = \beta_\pm$ modulo $\alpha^2 \mathcal{C}^\infty(\bar{X}_{\text{even}})$, where we note that r is a smooth function of μ , thus an *even* function of α , near $r = r_\pm$ in view of $\mu'(r_\pm) \neq 0$. The manifold \bar{X} is \bar{X}_{even} topologically, but with smooth functions of $\alpha = \mu^{1/2}$ added to the smooth structure. We denote $Y = \partial\bar{X} = \mathbb{S}^{n-2} \sqcup \mathbb{S}^{n-2}$.

By the analysis in Section 2, all zero resonant states u , written in the form (3.7) near Y , lie in the space $\mathcal{C}_{(0)}^\infty$, defined in (3.9). In the current setting, it is more natural to write differential forms as

$$u = u_{TT} + \alpha^{-1} dr \wedge u_{TN} + \alpha dt \wedge u_{NT} + \alpha dt \wedge \alpha^{-1} dr \wedge u_{NN}, \quad (4.3)$$

since $\alpha^{-1} dr$ has squared norm -1 (with respect to the metric g). We compute how the matching condition on the singular terms of u , encoded in the $\beta_\pm \alpha^{-1}$ entry of the matrix \mathcal{C} , changes when we thus change the basis of the form bundle: Namely, we have $\beta_\pm \alpha^{-1} d\alpha = (\mp 1 + \alpha^2 \mathcal{C}^\infty(\bar{X}_{\text{even}})) \alpha^{-1} \alpha^{-1} dr$; thus, for u written as in (4.3),

we have

$$u \in \mathcal{C}_{(0)}^\infty \iff \begin{pmatrix} u_{TT} \\ u_{TN} \\ u_{NT} \\ u_{NN} \end{pmatrix} \in \mathcal{C}_\pm \begin{pmatrix} \mathcal{C}^\infty(\overline{X}_{\text{even}}; \Lambda \mathbb{S}^{n-2}) \\ \mathcal{C}^\infty(\overline{X}_{\text{even}}; \Lambda \mathbb{S}^{n-2}) \\ \mathcal{C}^\infty(\overline{X}_{\text{even}}; \Lambda \mathbb{S}^{n-2}) \\ \mathcal{C}^\infty(\overline{X}_{\text{even}}; \Lambda \mathbb{S}^{n-2}) \end{pmatrix}$$

near $r = r_\pm$, where

$$\mathcal{C}_\pm = \begin{pmatrix} 1 & 0 & 0 & 0 \\ 0 & \alpha & \mp \alpha^{-1} & 0 \\ 0 & 0 & \alpha^{-1} & 0 \\ 0 & 0 & 0 & 1 \end{pmatrix}. \quad (4.4)$$

We now proceed to compute the explicit form of the operators d_p, δ_p and \square_p , where the subscript p indicates the form degree on which the operators act. First, we recall (3.2) and (3.4) in the form

$$d_p = \begin{pmatrix} d_{X,p} & 0 \\ \alpha^{-1} \partial_t & -\alpha^{-1} d_{X,p-1} \alpha \end{pmatrix}, \quad \delta_p = \begin{pmatrix} -\alpha^{-1} \delta_{X,p} \alpha & -\alpha^{-1} \partial_t \\ 0 & \delta_{X,p-1} \end{pmatrix},$$

and these operators act on forms $u = u_T + \alpha dt \wedge u_N$, with u_T and u_N differential forms on X . Writing forms on X as $v = v_T + \alpha^{-1} dr \wedge v_N$, we have

$$d_{X,p} = \begin{pmatrix} d_{\mathbb{S}^{n-2},p} & 0 \\ \alpha \partial_r & -d_{\mathbb{S}^{n-2},p-1} \end{pmatrix}. \quad (4.5)$$

In order to compute the codifferential, we observe that the volume density on X induced by h is given by $\alpha^{-1} r^{n-2} dr |d\omega|$, while the induced inner product on the fibers on the bundle of p -forms is

$$H_p = \begin{pmatrix} r^{-2p} \Omega_p & 0 \\ 0 & r^{-2(p-1)} \Omega_{p-1} \end{pmatrix},$$

where Ω_p is the fiber inner product on the p -form bundle on \mathbb{S}^{n-2} . Therefore,

$$\begin{aligned} \delta_{X,p} &= \begin{pmatrix} r^{-2} \delta_{\mathbb{S}^{n-2},p} & \partial_{r,p-1}^* \\ 0 & -r^{-2} \delta_{\mathbb{S}^{n-2},p-1} \end{pmatrix}, \\ \partial_{r,p-1}^* &= -\alpha r^{-(n-2)} r^{2(p-1)} \partial_r r^{-2(p-1)} r^{n-2}. \end{aligned} \quad (4.6)$$

We obtain:

Lemma 4.2. *In the bundle decomposition (4.3), we have*

$$d_p = \begin{pmatrix} d_{\mathbb{S}^{n-2},p} & 0 & 0 & 0 \\ \alpha \partial_r & -d_{\mathbb{S}^{n-2},p-1} & 0 & 0 \\ \alpha^{-1} \partial_t & 0 & -d_{\mathbb{S}^{n-2},p-1} & 0 \\ 0 & \alpha^{-1} \partial_t & -\partial_r \alpha & d_{\mathbb{S}^{n-2},p-2} \end{pmatrix} \quad (4.7)$$

and

$$\delta_p = \begin{pmatrix} -r^{-2} \delta_{\mathbb{S}^{n-2},p} & -\alpha^{-1} \partial_{r,p-1}^* \alpha & -\alpha^{-1} \partial_t & 0 \\ 0 & r^{-2} \delta_{\mathbb{S}^{n-2},p-1} & 0 & -\alpha^{-1} \partial_t \\ 0 & 0 & r^{-2} \delta_{\mathbb{S}^{n-2},p-1} & \partial_{r,p-2}^* \\ 0 & 0 & 0 & -r^{-2} \delta_{\mathbb{S}^{n-2},p-2} \end{pmatrix}. \quad (4.8)$$

Moreover,

$$\begin{aligned}
-r^2 \square_p = & \begin{pmatrix} \Delta_p & -2\alpha r d_{p-1} & 0 & 0 \\ -2\alpha r^{-1} \delta_p & \Delta_{p-1} & -r^2 \mu^{-1} \mu' \partial_t & 0 \\ 0 & -r^2 \mu^{-1} \mu' \partial_t & \Delta_{p-1} & -2\alpha r d_{p-2} \\ 0 & 0 & -2\alpha r^{-1} \delta_{p-1} & \Delta_{p-2} \end{pmatrix} \\
& + \begin{pmatrix} r^2 \alpha^{-1} \partial_{r,p}^* \alpha^2 \partial_r & 0 & 0 & 0 \\ 0 & r^2 \alpha \partial_r \alpha^{-1} \partial_{r,p-1}^* \alpha & 0 & 0 \\ 0 & 0 & r^2 \partial_{r,p-1}^* \partial_r \alpha & 0 \\ 0 & 0 & 0 & r^2 \partial_r \alpha \partial_{r,p-2}^* \end{pmatrix} \\
& + \begin{pmatrix} r^2 \mu^{-1} \partial_{tt} & 0 & 0 & 0 \\ 0 & r^2 \mu^{-1} \partial_{tt} & 0 & 0 \\ 0 & 0 & r^2 \mu^{-1} \partial_{tt} & 0 \\ 0 & 0 & 0 & r^2 \mu^{-1} \partial_{tt} \end{pmatrix}. \tag{4.9}
\end{aligned}$$

We can now compute the spaces \mathcal{K} and \mathcal{H} of zero resonances for \square and $d + \delta$ and deduce asymptotics for solutions of $(d + \delta)u = 0$:

Theorem 4.3. *On an n -dimensional Schwarzschild-de Sitter spacetime, $n \geq 4$, there exist two linearly independent 1-forms $u_{\pm} = f_{1,\pm}(r)\mu^{-1}dr + f_{2,\pm}(r)dt \in \mathcal{K}^1 = \ker \hat{\square}_1 \subset \mathcal{C}_{(0)}^{\infty}$,¹⁷ and we then have:*

$$\mathcal{K}^0 = \mathcal{H}^0 = \langle 1 \rangle, \quad \mathcal{K}^n = \mathcal{H}^n = \langle r^{n-2} dt \wedge dr \wedge \omega \rangle,$$

where ω denotes the volume form on the round sphere \mathbb{S}^{n-2} . Furthermore,

$$\begin{aligned}
\mathcal{K}^1 &= \langle u_+, u_- \rangle, \mathcal{H}^1 = 0, \quad \mathcal{K}^{n-1} = \langle \star u_+, \star u_- \rangle, \mathcal{H}^{n-1} = 0, \\
\mathcal{K}^k &= \mathcal{H}^k = 0, \quad k = 3, \dots, n-3.
\end{aligned}$$

For $n = 4$,

$$\mathcal{K}^2 = \mathcal{H}^2 = \langle \omega, r^{-2} dt \wedge dr \rangle,$$

while for $n > 4$,

$$\mathcal{K}^2 = \mathcal{H}^2 = \langle r^{-(n-2)} dt \wedge dr \rangle, \quad \mathcal{K}^{n-2} = \mathcal{H}^{n-2} = \langle \omega \rangle.$$

In particular, on 4-dimensional Schwarzschild-de Sitter space, if u is a solution of $(d + \delta)u = 0$ with smooth initial data, then the degree 0 component of u decays exponentially to a constant, the degree 1 and degree 3 components decay exponentially to 0, the degree 2 component decays exponentially to a linear combination of ω and $r^{-2} dt \wedge dr$, and the degree 4 component decays exponentially to a constant multiple of the volume form. Analogous statements hold on any n -dimensional Schwarzschild-de Sitter space, $n \geq 5$.

Proof. First, we observe that $H^k(\bar{X}) \cong H^k(\mathbb{S}^{n-2})$, $H^{k-1}(\bar{X}, \partial\bar{X}) \cong H^{n-k}(\bar{X}) \cong H^{n-k}(\mathbb{S}^{n-2})$ by Poincaré duality, and $H^{k-1}(\partial\bar{X}) \cong H^{k-1}(\mathbb{S}^{n-2}) \oplus H^{k-1}(\mathbb{S}^{n-2})$. Thus, the short exact sequence (3.52) immediately gives the dimensions of the spaces \mathcal{K}^k , and (3.51) gives the dimensions of \mathcal{H}^k for all values of k except $k = 1$ and $k = n - 1$.

We now compute \mathcal{H} and \mathcal{K} in the case $n = 4$. For $k = 0$, the short exact sequence (3.52) reads $0 \rightarrow H^0(\bar{X}) \oplus 0 \rightarrow \mathcal{K}^0 \rightarrow 0 \rightarrow 0$, and since $H^0(\bar{X}) = \langle [1] \rangle$, this suggests 1 as a resonant state for \square on 0-forms (i.e. functions), and indeed $\square 1 = 0$, hence

¹⁷The forms u_{\pm} have a simple explicit form, see (4.10) and Footnote 18.

$\mathcal{K}^0 = \langle 1 \rangle$. Theorem 3.20 also shows that $\mathcal{H}^0 = \mathcal{K}^0$. Then we immediately obtain $\mathcal{H}^4 = \mathcal{K}^4 = \langle \star 1 \rangle = \langle r^2 dt \wedge dr \wedge \omega \rangle$.

Next, we treat the form degree $k = 2$. Then (3.52) reads $0 \rightarrow H^2(\overline{X}) \oplus H^1(\overline{X}, \partial \overline{X}) \rightarrow \mathcal{K}^2 \rightarrow 0 \rightarrow 0$. Now $H^2(\overline{X}) = \langle [\omega] \rangle$, and a generator of $H^1(\overline{X}, \partial \overline{X})$ is given by the Poincaré dual of ω (which generates $H^2(\overline{X})$). This suggests the ansatz $u = f(r)\omega$ for an element of $\mathcal{K}^2 = \mathcal{H}^2$ (the latter equality following from (3.51)), and then $\star u$ will be the second element of a basis of \mathcal{K}^2 . Now, in the decomposition (4.3), we compute using Lemma 4.2 that $\widehat{\delta}_2(0)u = 0$ for $u = f(r)\omega$, and

$$\widehat{d}_2(0)u = \widehat{d}_2(0) \begin{pmatrix} f(r)\omega \\ 0 \\ 0 \\ 0 \end{pmatrix} = \begin{pmatrix} 0 \\ \alpha f'(r)\omega \\ 0 \\ 0 \end{pmatrix},$$

which vanishes precisely if $f(r)$ is constant.

The analysis of resonant states in form degree $k = 1$ is just a bit more involved. Since (3.52) now reads $0 \rightarrow 0 \oplus 0 \rightarrow \mathcal{K}^1 \rightarrow H^0(\mathbb{S}^2 \sqcup \mathbb{S}^2) \rightarrow 0$, every non-trivial element u of \mathcal{K}^1 fails to be in $L^2(\alpha|dh|)$, and in fact the singular behavior is expected to be $u = \mathcal{C}_\pm \tilde{u}$ with $\tilde{u}_{NT}|_{r=r_\pm} = c_\pm \in \mathbb{C}$, since $H^0(\mathbb{S}^2 \sqcup \mathbb{S}^2)$ is generated by locally constant functions, which are therefore constant on $r = r_-$ as well as on $r = r_+$. We thus make the ansatz

$$u = \alpha^{-1} f_1(r) \alpha^{-1} dr + \alpha dt \wedge \alpha^{-1} f_2(r). \quad (4.10)$$

We then compute

$$-\widehat{\square}_1(0)u = \begin{pmatrix} 0 \\ \alpha \partial_r \alpha^{-1} \partial_{r,0}^* f_1 \\ \partial_{r,0}^* \partial_r f_2 \\ 0 \end{pmatrix},$$

and by definition of $\partial_{r,p}^*$ in (4.6), this vanishes if and only if f_1 and f_2 satisfy the ODEs

$$\begin{aligned} \partial_r r^{-2} \partial_r r^2 f_1 &= 0, \\ r^{-2} \partial_r r^2 \partial_r f_2 &= 0. \end{aligned}$$

The general form of the solution is¹⁸

$$\begin{aligned} f_1(r) &= f_{11}r + f_{12}r^{-2}, \\ f_2(r) &= f_{21} + f_{22}r^{-1}, \end{aligned}$$

$f_{jk} \in \mathbb{C}$, $j, k = 1, 2$. Now recall that resonant states are elements of $\mathcal{C}_{(0)}^\infty$ and thus satisfy a matching condition in the singular components, which is captured by the matrix (4.4). Concretely, we require $f_2(r_-) = f_1(r_-)$ and $f_2(r_+) = -f_1(r_+)$; in terms of f_{jk} , $j, k = 1, 2$, these conditions translate into

$$\begin{pmatrix} r_- & r_-^{-2} & -1 & -r_-^{-1} \\ r_+ & r_+^{-2} & 1 & r_+^{-1} \end{pmatrix} \begin{pmatrix} f_{11} \\ f_{12} \\ f_{21} \\ f_{22} \end{pmatrix} = \begin{pmatrix} 0 \\ 0 \end{pmatrix}.$$

¹⁸On n -dimensional Schwarzschild-de Sitter space, the exponents 2 and -2 in these ODEs get replaced by $n-2$ and $2-n$, and the general forms of the solutions are $f_1(r) = f_{11}r + f_{12}r^{2-n}$ and $f_2(r) = f_{21} + f_{22}r^{3-n}$. The subsequent analysis of the matching conditions goes through with obvious modifications.

Since the 2×4 matrix on the left has rank 2, we get a 2-dimensional space of solutions. In fact, it is easy to see that we can freely specify the values $f_1(r_-)$ and $f_1(r_+)$, and f_1 and f_2 are then uniquely determined. To be specific, we can for instance define $u_+ \in \mathcal{K}^1$ to be the 1-form with $f_1(r_-) = 0$, $f_1(r_+) = 1$, and $u_- \in \mathcal{K}^1$ to be the 1-form with $f_1(r_-) = 1$, $f_1(r_+) = 0$, and we then have $\mathcal{K}^1 = \langle u_+, u_- \rangle$, as claimed.

Next, since $\mathcal{H}^1 \subset \mathcal{K}^1$, computing \mathcal{H}^1 simply amounts to finding all linear combinations of u_- and u_+ which are annihilated by both $\widehat{d}_1(0)$ and $\widehat{\delta}_1(0)$. But

$$\widehat{d}_1(0) \begin{pmatrix} 0 \\ \alpha^{-1} f_1(r) \\ \alpha^{-1} f_2(r) \\ 0 \end{pmatrix} = \begin{pmatrix} 0 \\ 0 \\ 0 \\ -\partial_r f_2 \end{pmatrix} = 0$$

requires f_2 to be constant, and

$$\widehat{\delta}_1(0) \begin{pmatrix} 0 \\ \alpha^{-1} f_1(r) \\ \alpha^{-1} f_2(r) \\ 0 \end{pmatrix} = \begin{pmatrix} -\alpha^{-1} \partial_{r,0}^* f_1 \\ 0 \\ 0 \\ 0 \end{pmatrix} = 0$$

implies $r^{-2} \partial_r r^2 f_1 = 0$, hence $f_1(r) = f_1(r_-)(r/r_-)^{-2}$. The matching condition requires $f_1(r_+) = f_1(r_-)(r_+/r_-)^{-2} = -f_2(r_+) = -f_2(r_-) = -f_1(r_-)$ and is therefore only satisfied if $f_1(r_-) = 0$, which implies $f_1 \equiv 0$ and $f_2 \equiv 0$. This shows that $\mathcal{H}^1 = 0$ and finishes the computation of the spaces of resonances for $n = 4$. The computation for spacetime dimensions $n \geq 5$ is completely analogous.

Finally, the statement about asymptotics of solutions to $(d + \delta)u = 0$ follows from the above computations combined with high energy estimates for $d + \delta$, which follow from those for \square , and Lemma 3.7.¹⁹ Once we check the normally hyperbolic nature of the trapping and show that the subprincipal symbol of \square (or a conjugated version thereof), relative to a *positive definite* fiber inner product, at the trapping is smaller than $\nu_{\min}/2$, where ν_{\min} is the minimal expansion rate in the normal direction at the trapped set, we can use Dyatlov's result [23] to obtain a spectral gap below the real line, i.e. the absence of resonances in a small strip below the reals, which combines with the general framework of [52] to yield the desired resonance expansion of solutions with exponentially decaying error terms; see [30, §§4.4, 5] for further details.

The dynamics of the Hamilton flow at the trapping only depend on properties of the scalar principal symbol g of \square . For easier comparison with [22, 52, 56], we consider the operator $\mathcal{P} = -r^2 \square$. We take the Fourier transform in $-t$, obtaining a family of operators on X depending on the dual variable τ , and then do a semiclassical rescaling, multiplying $\widehat{\mathcal{P}}$ by h^2 , giving a second order semiclassical differential operator P_h , with $h = |\tau|^{-1}$, and we then define $z = h\tau$. Introduce coordinates on T^*X by writing 1-forms as $\xi dr + \eta d\omega$, and let

$$\Delta_r := r^2 \mu = r^2(1 - \lambda r^2) - 2M_\bullet r^{5-n},$$

¹⁹If $(\widehat{d}(\sigma) + \widehat{\delta}(\sigma))^{-1}$ had a second order pole at 0, then solutions to $(d + \delta)u = 0$ would generically blow up linearly; the simplicity of the pole ensures that solutions stay bounded with the asymptotic stationary state given by an element of \mathcal{H} .

then the semiclassical principal symbol p of P_h is

$$p = \Delta_r \xi^2 - \frac{r^4}{\Delta_r} z^2 + |\eta|^2,$$

and correspondingly the Hamilton vector field is

$$H_p = 2\Delta_r \xi \partial_r - \left(\partial_r \Delta_r \xi^2 - \partial_r \left(\frac{r^4}{\Delta_r} \right) z^2 \right) \partial_\xi + H_{|\eta|^2}$$

We work with real z , hence $z = \pm 1$. We locate the trapped set: If $H_p r = 2\Delta_r \xi = 0$, then $\xi = 0$, in which case $H_p^2 r = 2\Delta_r H_p \xi = 2\Delta_r \partial_r (r^4/\Delta_r) z^2$. Recall the definition of the function $\tilde{\mu} = \mu/r^2 = \Delta_r/r^4$, then we can rewrite this as $H_p^2 r = -2\Delta_r \tilde{\mu}^{-2} (\partial_r \tilde{\mu}) z^2$. We have already seen that $\partial_r \tilde{\mu}$ has a single root $r_p \in (r_-, r_+)$, and $(r - r_p) \partial_r \tilde{\mu} < 0$ for $r \neq r_p$. Therefore, $H_p^2 r = 0$ implies (still assuming $H_p r = 0$) $r = r_p$. Thus, the only trapping occurs in the cotangent bundle over $r = r_p$: Indeed, define $F(r) = (r - r_p)^2$, then $H_p F = 2(r - r_p) H_p r$ and $H_p^2 F = 2(H_p r)^2 + 2(r - r_p) H_p^2 r$. Thus, if $H_p F = 0$, then either $r = r_p$, in which case $H_p^2 F = 2(H_p r)^2 > 0$ unless $H_p r = 0$, or $H_p r = 0$, in which case $H_p^2 F = 2(r - r_p) H_p^2 r > 0$ unless $r = r_p$. So $H_p F = 0, p = 0$ implies either $H_p^2 F > 0$ or $r = r_p, H_p r = 0$. Therefore, the trapped set in T^*X is given by

$$(r, \omega; \xi, \eta) \in \Gamma_h := \left\{ (r_p, \omega; 0, \eta) : \frac{r^4}{\Delta_r} z^2 = |\eta|^2 \right\},$$

and F is an escape function. The linearization of the H_p -flow at Γ_h in the normal coordinates $r - r_p$ and ξ equals

$$H_p \begin{pmatrix} r - r_p \\ \xi \end{pmatrix} = \begin{pmatrix} 0 & 2r_p^4 \tilde{\mu}|_{r=r_p} \\ 2(n-3)r_p^{-4}(\tilde{\mu}|_{r=r_p})^{-2} z^2 & 0 \end{pmatrix} \begin{pmatrix} r - r_p \\ \xi \end{pmatrix} + \mathcal{O}(|r - r_p|^2 + |\xi|^2),$$

where we used $\partial_{rr} \tilde{\mu}|_{r=r_p} = -2(n-3)r_p^{-4}$, which gives $\partial_r \tilde{\mu} = -2(n-3)r_p^{-4}(r - r_p) + \mathcal{O}(|r - r_p|^2)$. The eigenvalues of the linearization are therefore equal to $\pm \nu_{\min}$, where

$$\nu_{\min} = 2r_p \left(\frac{n-1}{1 - \frac{n-1}{n-3} r_p^2 \lambda} \right)^{1/2} > 0.$$

The expansion rate of the flow within the trapped set is 0 by spherical symmetry, since integral curves of H_p on Γ_h are simply unit speed geodesics of \mathbb{S}^{n-2} . This shows the normal hyperbolicity (in fact, r -normal hyperbolicity for every r) of the trapping.

It remains to bound the imaginary part of $\mathcal{P} = -r^2 \square_g$ in terms of ν_{\min} in order to obtain high energy estimates below the real line. More precisely, we need to show that

$$Q := |\tau|^{-1} \sigma_1 \left(\frac{1}{2i} (\mathcal{P} - \mathcal{P}^*) \right) < \frac{\nu_{\min}}{2}$$

at the trapped set (cf. the discussion in [30, §5.4]), where we take the adjoint with respect to some Riemannian inner product B , to be chosen, on the bundle $\Lambda^p \mathbb{S}^{n-2} \oplus \Lambda^{p-1} \mathbb{S}^{n-2} \oplus \Lambda^{p-1} \mathbb{S}^{n-2} \oplus \Lambda^{p-2} \mathbb{S}^{n-2}$; notice that Q is a self-adjoint section of the endomorphism bundle of this bundle. An obvious guess is to use $B = H \oplus H$ in the tangential-normal decomposition (3.6), thus

$$B = r^{-2p} \Omega_p \oplus r^{-2(p-1)} \Omega_{p-1} \oplus r^{-2(p-1)} \Omega_{p-1} \oplus r^{-2(p-2)} \Omega_{p-2}.$$

In this case, the expression (4.9) shows that the only parts of \mathcal{P} that are not symmetric with respect to B at the spacetime trapped set

$$\Gamma = \{(t, r_p, \omega; \tau, 0, \eta) : \frac{r^4}{\Delta_r} \tau^2 = |\eta|^2\},$$

are the (2, 3) and (3, 2) components; thus, taking adjoints with respect to B , we compute

$$Q = \begin{pmatrix} 0 & 0 & 0 & 0 \\ 0 & 0 & \pm r^2 \mu^{-1} \mu' & 0 \\ 0 & \pm r^2 \mu^{-1} \mu' & 0 & 0 \\ 0 & 0 & 0 & 0 \end{pmatrix}$$

at Γ , with the sign depending the sign of τ . Now $(\mu/r^2)' = 0$ at $r = r_p$ implies $\mu^{-1} \mu' r_p^2 = 2r_p$; the eigenvalues of Q are therefore $\pm 2r_p$, and they are bounded by $\nu_{\min}/2$ if and only if

$$r_p^2 \lambda > \frac{(5-n)(n-3)}{4(n-1)},$$

which in spacetime dimensions $n \geq 5$ is always satisfied. In dimension $n = 4$ however, the condition becomes $r_p^2 \lambda > 1/12$, or

$$9M_\bullet \Lambda > \frac{1}{4},$$

while the non-degeneracy condition (4.2) requires $9M_\bullet \Lambda < 1$. Therefore, only for very massive black holes or very large cosmological constants does the above choice of positive definite inner product B yield a sufficiently small imaginary part of \mathcal{P} . In fact, for parameters M_\bullet and Λ with $9M_\bullet \Lambda \leq 1/4$, the endomorphism Q is not bounded by $\nu_{\min}/2$ for any choice of B . To overcome this problem, one needs to allow B to be a *pseudodifferential inner product* on the form bundle, depending on the position in phase space rather than physical space; equivalently, one can replace \mathcal{P} by $\mathcal{Q}\mathcal{P}\mathcal{Q}^{-1}$, where \mathcal{Q} , an elliptic pseudodifferential operator acting on the form bundle, is chosen in such a way that $\mathcal{Q}\mathcal{P}\mathcal{Q}^{-1}$, relative to a Riemannian inner product on the form bundle, e.g. B , has (arbitrarily) small imaginary part, which is in particular bounded by $\nu_{\min}/2$. Pseudodifferential inner products are introduced by Hintz in a recent paper [28], where these claims are proved for general tensor-valued waves on Schwarzschild-de Sitter spacetimes with spacetime dimension $n \geq 4$; the point of view of pseudodifferential inner products shows precisely what the special structure of the subprincipal part of \square at the trapped set is that makes such a choice of a pseudodifferential inner product (equivalently, a choice of a conjugating operator \mathcal{Q}) possible. \square

We can in fact prove boundedness and asymptotics for solutions of the wave equation on differential forms in all form degrees as well. To begin, write

$$(\tilde{d}(\sigma) + \tilde{\delta}(\sigma))^{-1} = \sigma^{-1} A_{-1} + \mathcal{O}(1), \quad A_{-1} = \sum_{j=1}^4 \langle \cdot, \psi_j \rangle \phi_j, \quad (4.11)$$

near $\sigma = 0$, where $\{\phi_j\}_{j=1,\dots,4}$ is a basis of the space of resonant states and $\{\psi_j\}_{j=1,\dots,4}$ is a basis of the space $\mathcal{H}_* = \ker(\tilde{d}(0) + \tilde{\delta}(0))^*$ of dual states.²⁰ Therefore, we need to understand the dual states of $d + \delta$ in order to understand the

²⁰After choosing the ϕ_j , say, the ψ_j are uniquely determined.

order and structure of the pole of $\tilde{\square}(\sigma)^{-1} = ((\tilde{d}(\sigma) + \tilde{\delta}(\sigma))^{-1})^2$ as $\sigma = 0$. Notice here that the adjoint $(\tilde{d}(\sigma) + \tilde{\delta}(\sigma))^*$ acts on distributions on \tilde{X} which are *supported* at the Cauchy hypersurface $\partial\tilde{X}$ (see [31, Appendix B] for this and related notions). In particular, an element $\tilde{u} \in \ker(\tilde{d}(\sigma) + \tilde{\delta}(\sigma))^*$ satisfies $\tilde{u} \in \ker \tilde{\square}(\sigma)$ and is a supported distribution at $\partial\tilde{X}$, thus by local uniqueness, \tilde{u} vanishes in the hyperbolic region $\tilde{X} \setminus X$, hence $\text{supp } \tilde{u} \subset \bar{X}$.

Lemma 4.4. *The spaces \mathcal{H}_* and \mathcal{K}_* of dual states for $d + \delta$ and \square , respectively, on n -dimensional Schwarzschild-de Sitter space, $n \geq 4$, are graded by form degree, $\mathcal{H}_* = \bigoplus_{k=0}^n \mathcal{H}_*^k$, $\mathcal{K}_* = \bigoplus_{k=0}^n \mathcal{K}_*^k$, and have the following explicit descriptions:*

$$\begin{aligned} \mathcal{K}_*^0 &= \langle 1_X \rangle, \mathcal{H}_*^0 = 0, \quad \mathcal{K}_*^n = \langle 1_X r^{n-2} dt \wedge dr \wedge \omega \rangle, \mathcal{H}_*^n = 0, \\ \mathcal{H}_*^1 &= \langle \delta_{r=r_-} dr, \delta_{r=r_+} dr \rangle, \quad \mathcal{H}_*^{n-1} = \langle \delta_{r=r_-} dr \wedge \omega, \delta_{r=r_+} dr \wedge \omega \rangle, \\ \mathcal{H}_*^k &= 0, \quad k = 2, \dots, n-2, \end{aligned}$$

where ω denotes the volume form on the round sphere \mathbb{S}^{n-2} . Furthermore, $\mathcal{K}_*^1 = \mathcal{H}_*^1$, $\mathcal{K}_*^{n-1} = \mathcal{H}_*^{n-1}$ and

$$\mathcal{K}_*^k = 0, \quad k = 3, \dots, n-3.$$

For $n = 4$,

$$\mathcal{K}_*^2 = \langle 1_X \omega, 1_X r^{-2} dt \wedge dr \rangle,$$

while for $n > 4$,

$$\mathcal{K}_*^2 = \langle 1_X r^{2-n} dt \wedge dr \rangle, \quad \mathcal{K}_*^{n-2} = \langle 1_X \omega \rangle.$$

We have $\langle \phi, \psi \rangle = 0$ for all $\phi \in \mathcal{H}$, $\psi \in \mathcal{H}_*$.

Proof. For computing the dual resonant states, we need to compute the form of $\tilde{\square}(0)$ near the two components of $\partial\bar{X} = \mathbb{S}^{n-2} \sqcup \mathbb{S}^{n-2}$. Since dual states are supported in \bar{X}_{even} , it suffices to compute $\mathcal{C}_\pm^{-1} \tilde{\square}(0) \mathcal{C}_\pm$, since any smooth extension of this operator to \tilde{X} agrees with $\tilde{\square}(0)$ in X and to infinite order at $\partial\bar{X}_{\text{even}}$,²¹ thus the difference annihilates dual states. Using Lemma 4.2, we compute

$$\begin{aligned} -\mathcal{C}_\pm^{-1} \tilde{\square}_p(0) \mathcal{C}_\pm &= r^{-2} \begin{pmatrix} \Delta_p & 0 & 0 & 0 \\ 0 & \Delta_{p-1} & 0 & 0 \\ 0 & 0 & \Delta_{p-1} & 0 \\ 0 & 0 & 0 & \Delta_{p-2} \end{pmatrix} \\ &+ \begin{pmatrix} \alpha^{-1} \partial_{r,p}^* \alpha^2 \partial_r & -2\alpha^2 r^{-1} d_{p-1} & \pm 2r^{-1} d_{p-1} & 0 \\ -2r^{-3} \delta_p & \partial_r \alpha^{-1} \partial_{r,p-1}^* \alpha^2 & \pm(2(p-1) - (n-2))r^{-2} & \mp 2r^{-1} d_{p-2} \\ 0 & 0 & \alpha \partial_{r,p-1}^* \partial_r & -2\alpha^2 r^{-1} d_{p-2} \\ 0 & 0 & -2r^{-3} \delta_{p-1} & \partial_r \alpha \partial_{r,p-2}^* \end{pmatrix}, \end{aligned}$$

where the Laplace operators, differentials and codifferentials are the operators on \mathbb{S}^{n-2} . This does extend to an operator acting on smooth functions on $(r_\pm - \delta, r_\pm + \delta) \times \mathbb{S}^{n-2}$, $\delta > 0$ small, near r_\pm .

Now for $p = 0$, clearly $\alpha^{-1} \partial_{r,0}^* \alpha^2 \partial_r 1_X = \mp \alpha^{-1} \partial_{r,0}^* (\mu \delta_{r=r_\pm}) = 0$, hence $\mathcal{K}_*^0 = \langle 1_X \rangle$. (Observe that since $\tilde{\square}_0(0)$ is Fredholm of index 0 and has a 1-dimensional kernel according to Theorem 4.3, the space of dual 0-form resonances is 1-dimensional as well.) Likewise, for $p = n$, we have

$$\partial_r \alpha \partial_{r,n-2}^* (1_X r^{n-2} dt \wedge dr \wedge \omega) = -\partial_r \mu r^{n-2} \partial_r (1_X dt \wedge dr \wedge \omega) = 0,$$

²¹Since the Schwarzschild-de Sitter metric is analytic, we in fact do not have to make any choices.

confirming $\mathcal{K}_*^n = \langle 1_X r^{n-2} dt \wedge dr \wedge \omega \rangle$. By completely analogous arguments, we find $1_X r^{2-n} dt \wedge dr \in \mathcal{K}_*^2$ and $1_X \omega \in \mathcal{K}_*^{n-2}$.

In order to proceed, notice that $\tilde{d}(0) + \tilde{\delta}(0)$ maps \mathcal{K}_* into \mathcal{H}_* . Hence, we can find dual states for $d + \delta$ by applying $\tilde{d}(0) + \tilde{\delta}(0)$ to the dual states of \square that we have already identified. For this computation, we note

$$\begin{aligned} \mathcal{C}_\pm^{-1} \hat{d}_p(0) \mathcal{C}_\pm &= \begin{pmatrix} d_{\mathbb{S}^{n-2},p} & 0 & 0 & 0 \\ \partial_r & -d_{\mathbb{S}^{n-2},p-1} & 0 & 0 \\ 0 & 0 & -d_{\mathbb{S}^{n-2},p-1} & 0 \\ 0 & 0 & -\partial_r & d_{\mathbb{S}^{n-2},p-2} \end{pmatrix}, \\ \mathcal{C}_\pm^{-1} \hat{\delta}_p(0) \mathcal{C}_\pm &= \begin{pmatrix} -r^2 \delta_{\mathbb{S}^{n-2},p} & -\alpha^{-1} \partial_{r,p-1}^* \alpha^2 & \pm \alpha^{-1} \partial_{r,p-1}^* & 0 \\ 0 & r^{-2} \delta_{\mathbb{S}^{n-2},p-1} & 0 & \pm \alpha^{-1} \partial_{r,p-2}^* \\ 0 & 0 & r^{-2} \delta_{\mathbb{S}^{n-2},p-1} & \alpha \partial_{r,p-2}^* \\ 0 & 0 & 0 & -r^{-2} \delta_{\mathbb{S}^{n-2},p-2} \end{pmatrix}. \end{aligned}$$

Thus, $(\tilde{d}_0(0) + \tilde{\delta}_0(0))1_X$ and $(\tilde{d}_2(0) + \tilde{\delta}_2(0))(1_X r^{2-n} dt \wedge dr)$ are both linear combinations of $\delta_{r=r_\pm} dr$, hence $\delta_{r=r_\pm} dr \in \mathcal{H}_*^1 \subset \mathcal{K}_*^1$, and similarly $(\tilde{d}_n(0) + \tilde{\delta}_n(0))(1_X \star 1)$ and $(\tilde{d}_{n-2}(0) + \tilde{\delta}_{n-2}(0))(1_X \omega)$ are both linear combinations of $\delta_{r=r_\pm} dr \wedge \omega$, hence $\delta_{r=r_\pm} dr \wedge \omega \in \mathcal{H}_*^{n-1} \subset \mathcal{K}_*^{n-1}$.

We have therefore identified 4 and 8 linearly independent dual states for $d + \delta$ and \square , which is equal to the dimensions of \mathcal{H} and \mathcal{K} , respectively, and since $\tilde{d}(0) + \tilde{\delta}(0)$ and $\tilde{\square}(0)$ have index 0, all dual states are linear combinations of these, i.e. we have thus identified a basis of the spaces of dual states. The orthogonality of resonant and dual states for $d + \delta$ follows immediately from the explicit forms of both derived in Theorem 4.3 and in this lemma: All dual states have form degree 1 or $n - 1$, while all resonant states have form degree 0, 2, $n - 2$ or n . \square

The orthogonality statement in Lemma 4.4 combined with (4.11) immediately gives $A_{-1}^2 = 0$, hence the coefficient of σ^{-2} in the Laurent expansion of $\hat{\square}(\sigma)^{-1}$ at $\sigma = 0$ vanishes. For precisely those form degrees $0 \leq p \leq n$ for which \mathcal{K}^p is non-trivial, $\hat{\square}(\sigma)^{-1}$ does have a simple pole at $\sigma = 0$, and

$$\hat{\square}_p(\sigma)^{-1} = \sigma^{-1} \sum_{j=1}^{\dim \mathcal{K}^p} \langle \cdot, \psi'_j \rangle \phi'_j + \mathcal{O}(1),$$

where ϕ'_j and ψ'_j run over a basis of $\ker \hat{\square}_p(0) \cong \mathcal{K}^p$ and $\mathcal{K}_*^p = \ker \hat{\square}_p(0)^*$, respectively.²²

Theorem 4.5. *On 4-dimensional Schwarzschild-de Sitter space, if $0 \leq p \leq 4$ and u is a differential form of degree p which solves $\square u = 0$ with smooth initial data, then u decays exponentially to*

- a constant for $p = 0$,
- a linear combination of u_+ and u_- , defined in the statement of Theorem 4.3, for $p = 1$,
- a linear combination of ω and $r^{-2} dt \wedge dr$ for $p = 2$,
- a linear combination of $\star u_+$ and $\star u_-$ for $p = 3$ and
- a constant multiple of $r^2 dt \wedge dr \wedge \omega$ for $p = 4$.

²²After choosing the ϕ'_j , the ψ'_j are uniquely determined, and vice versa.

Analogous statements hold on any n -dimensional Schwarzschild-de Sitter space, $n \geq 5$.

5. RESULTS FOR KERR-DE SITTER SPACE

We now prove that some of the results obtained in the previous section for the 4-dimensional Schwarzschild-de Sitter spacetime are stable under perturbations, which allows us to draw conclusions about asymptotics for solutions of $(d + \delta)u = 0$ or $\square u = 0$ on Kerr-de Sitter space with very little effort. Thus, fixing the black hole mass M_\bullet and the cosmological constant $\Lambda > 0$, denote by g_a the Kerr-de Sitter metric with angular momentum a ; thus, g_0 is the Schwarzschild-de Sitter metric.²³ Only very basic facts about the metric will be used; we refer the reader to [52, §6] for details and further information. We will write δ_{g_a} for the codifferential with respect to the metric g_a . We furthermore denote by $M = \mathbb{R}_t \times X$ the domain of exterior communications, and by $\tilde{M} = \mathbb{R}_{t_*} \times \tilde{X}$ the ‘extended’ spacetime.

To begin, recall that the scalar wave equation (and by essentially the same arguments the wave equation on differential forms, since the principal symbol of the Hodge d’Alembertian is scalar, see also [51, §4] for a discussion in a related context) on the Kerr-de Sitter spacetime fits into the microlocal framework developed in [52]. In particular, asymptotics for waves follow directly from properties of the Mellin transformed normal operator family, and moreover the analysis of the latter is stable under perturbations. In the present context, this concretely means that for any $\epsilon > 0$, there exists $a_\epsilon > 0$ such that for all angular momenta a with $|a| < a_\epsilon$, the meromorphic family of operators $\mathcal{R}_a(\sigma) := (\tilde{d}(\sigma) + \widetilde{\delta_{g_a}}(\sigma))^{-1}$ has no poles in $|\sigma| \geq \epsilon$, $\text{Im } \sigma \geq 0$, and such that moreover all poles in $|\sigma| < \epsilon$ are perturbations of the pole of $\mathcal{R}_0(\sigma)$ at 0, in the sense that all poles of $\mathcal{R}_a(\sigma)$ in $|\sigma| < \epsilon$ are simple, and the sum of the dimensions of the resonant spaces at each resonance (i.e. pole of $\mathcal{R}_a(\sigma)$) equals the corresponding sum for the Schwarzschild-de Sitter metric, which equals 4 by Theorem 4.3. Now, Lemma 4.4 suggests considering dual resonant states instead (which have a simpler form); the same stability result as for $\mathcal{R}_a(\sigma)$ holds for $\mathcal{R}_a^*(\sigma) := ((\tilde{d}(\sigma) + \widetilde{\delta_{g_a}}(\sigma))^*)^{-1}$. However, just as in the case of Schwarzschild-de Sitter space, we can immediately write down 4 linearly independent dual 0-resonant states for $d + \delta_{g_a}$: Namely, apply $\tilde{d}(0) + \widetilde{\delta_{g_a}}(0)$ to 1_X (this is a dual resonant state for \square_{g_a}), which produces a sum of δ -distributions supported at the horizons $r = r_\pm$, and splitting this up into the part supported at r_- and the part supported at r_+ , we obtain 2 linearly independent dual resonant states for $d + \delta$ in form degree 1. The same procedure can be applied to $\star_{g_a} 1_X$, yielding 2 linearly independent dual resonant states for $d + \delta$ in form degree 3 (which are simply the Hodge duals of the dual states in form degree 1). Hence,

$$\mathcal{H}_a := \ker(\tilde{d}(0) + \widetilde{\delta_{g_a}}(0)), \quad (5.1)$$

which has the same dimension as

$$\mathcal{H}_{a,*} := \ker(\tilde{d}(0) + \widetilde{\delta_{g_a}}(0))^*, \quad (5.2)$$

²³Assuming the non-degeneracy condition (4.2), which ensures that the cosmological horizon lies outside the black hole event horizon, the same will be true for small $|a|$, which is the setting in which work here. In general, one would need to assume that Λ , M_\bullet and a are such that the non-degeneracy condition [52, (6.2)] holds.

is at least 4-dimensional for small $|a|$, but it is also at most 4-dimensional by the above perturbation stability argument! Hence, for small $|a|$, we deduce that 0 is the only pole of $\mathcal{R}_a(\sigma)$, i.e. the only resonance of $d + \delta_{g_a}$, in $\text{Im } \sigma \geq 0$ (and also the only pole of $\mathcal{R}_a^*(\sigma)$ in this half space), and is simple.

We can use this in turn to prove the stability of the zero resonance for \square_{g_a} in all form degrees. Let $\pi_k: \mathcal{C}^\infty(\widetilde{M}; \Lambda \widetilde{M}) \rightarrow \mathcal{C}^\infty(\widetilde{M}; \Lambda \widetilde{M})$ denote the projection onto differential forms with pure form degree $k \in \{0, \dots, 4\}$, which induces a map on $\mathcal{C}^\infty(\widetilde{X}; \Lambda \widetilde{X} \oplus \Lambda \widetilde{X})$. Let

$$\mathcal{K}_a := \ker \widetilde{\square_{g_a}}(0) = \bigoplus_{k=0}^4 \mathcal{K}_a^k \quad (5.3)$$

be the grading of the zero resonant space of \square_{g_a} by form degree, likewise

$$\mathcal{K}_{a,*} := \ker \widetilde{\square_{g_a}}(0)^* = \bigoplus_{k=0}^4 \mathcal{K}_{a,*}^k \quad (5.4)$$

for the space of dual resonant states. Observe that $\pi_k \mathcal{H}_a \subseteq \mathcal{K}_a^k$, since $u \in \mathcal{H}_a$ implies $0 = \pi_k \square_{g_a} u = \square_{g_a} \pi_k u$. Now, since $\square_{g_a} 1 = 0$, we have $\mathcal{K}_a^0 = \langle 1 \rangle$ for small $|a|$ by stability, likewise $\mathcal{K}_a^4 = \langle \star_{g_a} 1 \rangle$. Furthermore, \mathcal{K}_a^2 is at most 2-dimensional for small $|a|$ (since \mathcal{K}_0^2 is 2-dimensional), but also $\mathcal{K}_a^2 \supseteq \pi_2 \mathcal{H}_a$; now $\pi_2 \mathcal{H}_0$ is 2-dimensional by Theorem 4.3 and \mathcal{H}_a depends smoothly on a , thus $\mathcal{K}_a^2 = \pi_2 \mathcal{H}_a$ is 2-dimensional for small $|a|$; therefore $\mathcal{K}_a^2 = \pi_2 \mathcal{H}_a$ is 2-dimensional. Finally, we have $\mathcal{H}_{a,*}^1 \subseteq \mathcal{K}_{a,*}^1$, hence by the analysis of $d + \delta_{g_a}$ above, $\mathcal{K}_{a,*}^1$, hence \mathcal{K}_a^1 , is at least 2-dimensional, but since \mathcal{K}_0^1 is 2-dimensional, we must in fact have $\dim \mathcal{K}_a^1 = 2$ for small $|a|$; likewise $\dim \mathcal{K}_a^3 = 2$. Hence, we have $\dim \mathcal{K}_a^k = \dim \mathcal{K}_0^k$ for $k = 0, \dots, 4$, which in particular means that the zero resonance of \square_{g_a} is the only resonance in $\text{Im } \sigma \geq 0$, and the resonance is simple.

We now summarize the above discussion, including a small improvement. The following theorem is completely parallel to Theorem 4.3, Lemma 4.4 and Theorem 4.5 for Schwarzschild-de Sitter spacetimes, extending these to Kerr-de Sitter spacetimes with small angular momentum:

Theorem 5.1. *For small $|a|$, the only resonance of $d + \delta_{g_a}$ in $\text{Im } \sigma \geq 0$ is a simple resonance at $\sigma = 0$, likewise for \square_{g_a} . The spaces \mathcal{H}_a and $\mathcal{H}_{a,*}$ of resonant and dual resonant states for $d + \delta_{g_a}$ are graded by form degree as $\mathcal{H}_a = \bigoplus_{k=0}^4 \mathcal{H}_a^k$, $\mathcal{H}_{a,*} = \bigoplus_{k=0}^4 \mathcal{H}_{a,*}^k$, in particular $\mathcal{H}_a^k = \ker \widetilde{d}_k(0) \cap \ker (\widetilde{\delta_{g_a}})_k(0)$,²⁴ with*

$$\mathcal{H}_a^0 = \langle 1 \rangle, \quad \mathcal{H}_a^1 = 0, \quad \mathcal{H}_a^2 = \langle u_{a,1}, u_{a,2} \rangle, \quad \mathcal{H}_a^3 = 0, \quad \mathcal{H}_a^4 = \langle \star_{g_a} 1 \rangle$$

for some 2-forms $u_{a,1}, u_{a,2}$, which can be chosen to depend smoothly on a ,²⁵ with $u_{0,1} = r^{-2} dt \wedge dr$, $u_{0,2} = \omega$ in the notation of Theorem 4.3, and

$$\begin{aligned} \mathcal{H}_{a,*}^0 &= 0, & \mathcal{H}_{a,*}^1 &= \langle \delta_{r=r_-} dr, \delta_{r=r_+} dr \rangle, \\ \mathcal{H}_{a,*}^2 &= 0, & \mathcal{H}_{a,*}^3 &= \star_{g_a} \mathcal{H}_{a,*}^1, & \mathcal{H}_{a,*}^4 &= 0. \end{aligned}$$

For the spaces \mathcal{K}_a and $\mathcal{K}_{a,*}$ of resonant and dual resonant states for \square_{g_a} , we have

$$\mathcal{K}_a^0 = \mathcal{H}_a^0, \quad \mathcal{K}_a^1 = \langle u_{a,+}, u_{a,-} \rangle, \quad \mathcal{K}_a^2 = \mathcal{H}_a^2, \quad \mathcal{K}_a^3 = \star_{g_a} \mathcal{K}_a^1, \quad \mathcal{K}_a^4 = \mathcal{H}_a^4$$

²⁴The subscript denotes the degree of differential forms on which the respective operator acts.

²⁵We derive an explicit expression in Remark 5.4 below.

for some 1-forms $u_{a,\pm}$, which can be chosen to depend smoothly on a , with $u_{0,\pm} = u_{\pm}$ in the notation of Theorem 4.5, and

$$\begin{aligned}\mathcal{K}_{a,*}^0 &= \langle 1_X \rangle, & \mathcal{K}_{a,*}^1 &= \mathcal{H}_{a,*}^1, \\ \mathcal{K}_{a,*}^2 &= \langle 1_X u_{a,1}, 1_X u_{a,2} \rangle, & \mathcal{K}_{a,*}^3 &= \mathcal{H}_{a,*}^3, & \mathcal{K}_{a,*}^4 &= \langle \star_{g_a} 1_X \rangle.\end{aligned}$$

In particular, the form degree k part of a solution u to $(d + \delta_{g_a})u = 0$, resp. $\square_{g_a} u = 0$, with smooth initial data decays exponentially to an element of \mathcal{H}_a^k , resp. \mathcal{K}_a^k , for $k = 0, \dots, 4$.

Remark 5.2. Since for all $k = 0, \dots, 4$, either $\mathcal{H}_a^k = 0$ or $\mathcal{H}_{a,*}^k = 0$, hence \mathcal{H}_a and \mathcal{H}_* are orthogonal, we obtain another proof, as in the Schwarzschild-de Sitter case, of the fact that \square_{g_a} acting on differential forms only has a simple resonance at 0.

Proof of Theorem 5.1. We only need to prove that the space \mathcal{H}_a is graded by form degree: Let $\pi_{\text{even}} = \pi_0 + \pi_2 + \pi_4$ denote the projection onto even form degree parts, then since $d + \delta_{g_a}$ maps even degree forms to odd degree forms and vice versa, π_{even} maps \mathcal{H}_a into itself. Now suppose $u \in \pi_{\text{even}} \mathcal{H}_a$, and write $u = u_0 + u_2 + u_4$ with $u_k = \pi_k u$, $k = 0, 2, 4$. Then $0 = \pi_1(d + \delta_{g_a})u = du_0 + \delta_{g_a} u_2$,²⁶ and applying δ_{g_a} to this equation gives $0 = \square_{g_a} u_0$, which implies $u_0 \in \mathcal{K}_a^0$, i.e. u_0 is a constant, as discussed before the statement of the theorem. Likewise, $u_4 \in \mathcal{K}_a^4$, so u_4 is the Hodge dual of a constant. Therefore, $d + \delta_{g_a}$ annihilates both u_0 and u_4 , hence $u_2 \in \mathcal{H}_a$. This argument shows that in fact $\pi_2 \mathcal{H}_a \subset \mathcal{H}_a$. Since $\pi_2 \mathcal{H}_a$ is 2-dimensional, as noted above, we have

$$\langle 1 \rangle \oplus \pi_2 \mathcal{H}_a \oplus \langle \star_{g_a} 1 \rangle \subseteq \mathcal{H}_a,$$

with both sides having the same dimension (namely, 4), and thus equality holds, providing the grading of \mathcal{H}_a by form degree. \square

This in particular proves Theorem 1.

Remark 5.3. Observe that *all* ingredients in the Fredholm analysis of the normal operator family, which here in particular involves estimates at normally hyperbolic trapping, as well as *all* of the above arguments which lead to a characterization of the spaces of resonances are stable in the sense that they apply to *any* stationary perturbation of a given Schwarzschild-de Sitter spacetime (4-dimensional for the above, but similar arguments apply in all spacetime dimensions ≥ 4), not only to slowly rotating Kerr-de Sitter black holes.

In fact, using the analysis of operators with non-smooth coefficients developed in [27] and extended in [30], we can deduce decay and expansions in the exact same form as in the above theorem for waves on spacetimes which are merely ‘asymptotically stationary’ and close to Schwarzschild-de Sitter, i.e. for which the metric tensor differs from a stationary metric close to Schwarzschild-de Sitter by an exponentially decaying symmetric 2-tensor (with suitable regularity). This shows at once that *quasilinear* wave equations on differential forms of the form $\square_{g(u, \nabla u)} u = q(u, \nabla u)$ with small initial data can be solved globally, provided $g(0, 0)$ is close to the Schwarzschild-de Sitter metric, and the non-linearity q annihilates 0-resonant states; to give an (artificial) example, on 2-forms, one could take $q(u, \nabla u) = |du|^2 u$.

²⁶We use the identification of resonant states with t_* -independent forms as in the proof of Theorem 3.20.

Remark 5.4. In the case of the Kerr-de Sitter metric, we can in fact explicitly write down $u_{a,1} \in \mathcal{H}_a^2$ (and then take $u_{a,2} = \star_{g_a} u_{a,1}$ to obtain a basis of \mathcal{H}_a^2). Indeed, on the Kerr spacetime, Andersson and Blue [2] give the values of the spin coefficients of the Maxwell field for the Coulomb solution in [2, §3.1], and reconstructing the Maxwell field itself (in the basis given by wedge products of differentials of the Boyer-Lindquist coordinates t, r, θ, ϕ) is then an easy computation using the explicit form of the null tetrad given in [2, Introduction, §2.4].²⁷ A tedious but straightforward calculation shows that the resulting 2-form

$$u_{a,1} := F_{a,TR}(r, \theta) (dt - a \sin^2 \theta d\phi) \wedge dr \\ + F_{a,\Theta\Phi}(r, \theta) \sin \theta d\theta \wedge (a dt - (r^2 + a^2) d\phi)$$

with

$$F_{a,TR}(r, \theta) = \frac{r^2 - a^2 \cos^2 \theta}{(r^2 + a^2 \cos^2 \theta)^2}, \quad F_{a,\Theta\Phi}(r, \theta) = \frac{2ar \cos \theta}{(r^2 + a^2 \cos^2 \theta)^2}$$

is a solution of Maxwell's equations on Kerr-de Sitter space as well, i.e. when the cosmological constant is positive.

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²⁷In the definition of ϕ_0 in [2, §2.4], the second summand $F[\hat{\Theta}, \hat{\Phi}_{\text{PNV}}]$ should be replaced by $F[\hat{\mathbf{m}}, \hat{\mathbf{m}}]$ to yield the correct result, see also [8, Equation (2)†].

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DEPARTMENT OF MATHEMATICS, STANFORD UNIVERSITY, CA 94305-2125, USA

E-mail address: `phintz@math.stanford.edu`

E-mail address: `andras@math.stanford.edu`